

Analysis of $B_c \rightarrow D_s^* \ell^+ \ell^-$ in the Standard Model Beyond Third Generation

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We study the FCNC $B_c \rightarrow D_s^* \ell^+ \ell^-$ ($\ell = \mu, \tau$) transition in the Standard Model with fourth generation (SM4). Taking fourth generation quark mass $m_{t'}$ of about 300 to 600 GeV with the CKM matrix elements $|V_{t'b}^* V_{ts}|$ in the range $(0.03 - 1.2) \times 10^{-2}$ and using the new CP odd phase (ϕ_{sb}) to be 90° , the analysis of decay rates, forward-backward asymmetries (FBA), lepton polarization asymmetries ($P_{L,N,T}$) and the helicity fractions of D_s^* meson ($f_{L,T}$) in $B_c \rightarrow D_s^* \ell^+ \ell^-$ ($\ell = \mu, \tau$) is made. It is found that in the fourth generation parameter space the above mentioned physical observables deviates sizably from their SM values in $B_c \rightarrow D_s^* \mu^+ \mu^-$. Furthermore, an optimum shift in the zero position of the FBA in this decay has also been pointed out. Compared to the dimuon case, the SM4 effects are somewhat mild in the decay rate and in the FBA for $B_c \rightarrow D_s^* \tau^+ \tau^-$ decay. However, they came up in a distinguishing way in longitudinal and transverse lepton polarizations and also in the helicity fractions of the D_s^* meson which differ distinctively from their SM values. Thus the study of these physical observables will provide us useful information to probe new physics and helps us to search the fourth generation of quarks (t', b') in an indirect way.

I. INTRODUCTION

It is well known that the Standard Model (SM) includes three generations of fermions, but it does not prohibit the fourth generation. The restrictions on the number of fermion generations come from the QCD asymptotic freedom which constraint them to nine. Therefore, shortly after the measurement of the third generation, a fourth generation was an obvious extension.

Interest in the fourth generation Standard Model (SM4) was fairly high in the 1980s until the electroweak precision data seemed to rule it out. The other reason which shook the interest in the fourth generation was the measurement of the number of light neutrinos at the Z pole that showed only three light neutrinos could exist. However, the discovery of neutrino oscillations suggested the possibility of a mass scale beyond the SM, and the models with the sufficiently massive neutrino became acceptable [1]. Though the early study of the EW precision measurements ruled out a fourth generation [2], however it was subsequently pointed out [3] that if the fourth generation masses are not degenerate, then the EW precision data do not prohibit the fourth generation [4]. Therefore, the SM can be simply extended with a sequential repetition as four quark and four lepton left handed doublets and corresponding right

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handed singlets.

The possible sequential fourth generation may play an important role in understanding the well known problem of CP violation and flavor structure of standard theory [5–11], electroweak symmetry breaking [12–15], hierarchies of fermion mass and mixing angle in quark/lepton sectors [16–20]. A thorough discussion on the theoretical and experimental aspects of the fourth generation can be found in ref. [21].

On the experimental side, recent searches by the CDF collaboration for direct production of fourth generation up-type quark (t') and down-type quark (b') found $m_{t'} > 335$ GeV [22] and $m_{b'} > 385$ GeV [23], assuming $Br(t' \rightarrow Wq, (q = d, s, b)) = 100\%$ and $Br(b' \rightarrow Wt) = 100\%$ respectively. This indeed suggest that the fourth generation fermion must be heavy which supports the scenario of compositeness. The underlying assumption to perform these searches is that $m_{t'} - m_{b'} < M_W$ and negligible mixing of the (t', b') states with the two lightest quark generations. To account for EW precision data such conditions are generally required for the SM4 with the one Higgs doublet [24]. Moreover, when a fourth generation of fermions is embedded in theories beyond the SM, the large splitting case ($m_{t'} - m_{b'} > M_W$) and the inverted scenario ($m_{t'} < m_{b'}$) have not been excluded. Recently, it has also been shown [25] that the precision EW data can accommodate ($m_{t'} - m_{b'} > M_W$) if there are two Higgs doublets. Thus there is no uniquely interesting set of assumptions under which experimental data must be interpreted [26] and the determination of allowed parameter space of fourth generation fermions will be an important goal of the LHC era. The large values of the masses of fourth generation would provide special advantage to new interactions originating at a higher scale and the precise determination of the fourth generation quark properties may present the existence of physics beyond the SM.

It is necessary to mention here that these new particles are heavy in nature, consequently they are hard to produce in the accelerators. Therefore, we have to go for some alternate scenarios where we can find their influence. In this regard, the Flavor Changing Neutral Current (FCNC) transitions provide an ideal platform to establish new physics (NP). This is because of the fact that FCNC transitions are not allowed at tree level in the SM and are allowed at loop level through GIM mechanism which can get contributions of NP from newly proposed particles via loop diagrams. Among different FCNC transitions the one $b \rightarrow s$ transition plays a pivotal role to perform efficient tests of NP scenarios [27–35]. It is also the fact that CP violation in $b \rightarrow s$ transitions is predicted to be very small in the SM, thus, any experimental evidence for sizable CP violating effects in the B system would clearly point towards a NP scenario. However, among the other NP scenarios such as the Littlest Higgs model with T-parity (LHT) and Randall-Sundrum (RS) models, the study of FCNC transitions in SM4 contain much fewer parameters and has the possibility of having simultaneously sizable NP effects in the K and B systems compared to the above mentioned NP models. In this context the constraint on the mixing between the fourth and third generation by using FCNC processes have been studied [36] along with the effects of the sequential fourth generation on different physical observables in $B_{(d,s)}$, K and D decays, see ref. [37] for an incomplete list. The study of the B system will be even more complete if one consider the similar decays of the charmed B mesons (B_c). As B_c is a bound state of two heavy quarks b and c , therefore it is rich in phenomenology compared to other B mesons. In literature, some of the possible radiative and semileptonic exclusive decays of B_c mesons like $B_c \rightarrow (\rho, K^*, D_s^*, B_u^*)\gamma$, $B_c \rightarrow \ell\nu\gamma$, $B_c \rightarrow B_u^*\ell^+\ell^-$, $B_c \rightarrow D_1^0\ell\nu$, $B_c \rightarrow D_{s0}^*\ell^+\ell^-$ and $B_c \rightarrow D_{s,d}^*\ell^+\ell^-$ have been studied using the frame work of relativistic constituent quark model [39], QCD Sum Rules and Light Cone Sum Rules [38]. In this work we will focus on the semileptonic $B_c \rightarrow D_s^*\ell^+\ell^-$ decays in SM4.

The special feature of the semileptonic $B_c \rightarrow D_s^*\ell^+\ell^-$ decays compared to the other B meson decays such as

$B^0 \rightarrow (K^*, K_1, \rho, \pi)\ell^+\ell^-$ is that this decay can occur in two distinct ways i.e. through FCNC transitions and through Weak Annihilations (WA). In the ordinary charged B meson decays the WA contributions are very small and can be safely ignored. However, for B_c meson these WA contributions are proportional to the CKM matrix elements $V_{cb}V_{cs}^*$ and hence can not be ignored. *The decay under discussion here has already been studied in the Universal Extra Dimension (UED) model [41] where it has been seen that the WA contribution suppress the NP effects coming through UED model in different physical observables. Therefore, it is interesting to see how SM4 effects in different physical observables will behave in the presence of these WA contributions.*

While working on the exclusive B -meson decays the main job is to calculate the form factors which are the non perturbative quantities and are the scalar functions of the square of momentum transfer (q^2). In literature the form factors for $B_c \rightarrow D_s^*\ell^+\ell^-$ ($\ell = \mu, \tau$) decays were calculated using different approaches, such as light front constituent quark models, a relativistic quark model, QCD sum rules and through Ward identities [39–44]. In this work we will use QCD Sum rules form factors [43, 44] to study the fourth generation effects on different physical observables such as branching ratio, forward-backward asymmetry, polarization asymmetries of leptons and the helicity fractions of final state meson (D_s^*) for these decays. As the expected number of events for the production of B_c meson at the Large Hadron Collider (LHC) are about $10^8 - 10^{10}$ per year [42] therefore, we hope that the phenomenological study of the B_c meson could provide us a valuable tool for distinguishing the SM4 effects from the SM and other NP scenarios.

The paper is organized as follows. In Sec. II, we discuss the two different contributions to the amplitude of $B_c \rightarrow D_s^*\ell^+\ell^-$ decays which were named as WA and penguin contributions. These can be parameterized in terms of the form factors, where the values of the form factors appearing in the calculation of WA amplitude will be used from [43] and for the penguin contributions will be taken from [44]. Section III presents the basic formulas for physical observables like decay rate, lepton forward-backward asymmetry, helicity fractions of D_s^* meson and the polarization asymmetries of the final state lepton. The numerical study of these observables will be given in Section V and the Section VI gives the summary of the main outcomes of this study.

II. THEORETICAL FRAMEWORK FOR $B_c \rightarrow D_s^*\ell^+\ell^-$ DECAYS

A. Weak Annihilation Amplitude

The charmed B -meson (B_c) is made up of two different heavy flavors, b -quark and c -quark, which brings WA contributions into the play. Following the procedure given in refs. [45, 46] for $B_c \rightarrow D_s^*\gamma$ decay the WA amplitude for the decay $B_c \rightarrow D_s^*\ell^+\ell^-$ can be written as

$$\mathcal{M}^{\text{WA}} = \frac{G_F \alpha}{2\sqrt{2}\pi} \frac{f_{D_s^*} f_{B_c}}{q^2} V_{cb} V_{cs}^* \left[-i\epsilon_{\mu\nu\alpha\beta} \varepsilon^{*\nu} p^\alpha q^\beta F_V^{D_s^*}(q^2) + (\varepsilon \cdot qp_\mu + p \cdot q\varepsilon_\mu) F_A^{D_s^*}(q^2) \right] \bar{l}\gamma^\mu l \quad (1)$$

where f_{B_c} and $f_{D_s^*}$ are the decay constants of B_c and D_s^* mesons, respectively. The functions $F_V^{D_s^*}(q^2)$ and $F_A^{D_s^*}(q^2)$ are the weak annihilation form factors which are calculated in QCD Sum Rules [43] and can be parameterized as:

$$F_{V,A}^{D_s^*}(q^2) = \frac{F_{V,A}^{D_s^*}(0)}{1 + \alpha\hat{q} + \beta\hat{q}^2} \quad (2)$$

where $\hat{q} = q^2/M_{B_c}^2$. The values of the form factors $F_V^{D_s^*}(0)$ and $F_A^{D_s^*}(0)$ along with the given values of parameters α and β were calculated by Azizi et al. [43] which are summarized in Table I.

TABLE I: $B_c \rightarrow D_s^*$ form factors corresponding to WA in the QCD Sum Rules. $F(0)$ denotes the value of form factors at $q^2 = 0$ while α and β are the parameters in the parameterizations shown in Eq. (16)[43].

$F(q^2)$	$F(0)$	α	β
$F_V^{D_s^*}(q^2)$	0.23	-1.25	-0.097
$F_A^{D_s^*}(q^2)$	0.25	-0.10	-0.097

B. Penguin Amplitude

At quark level the decay $B_c \rightarrow D_s^* \ell^+ \ell^-$ ($\ell = \mu, \tau$) is governed by the transition $b \rightarrow s \ell^+ \ell^-$ for which the effective Hamiltonian can be written as

$$H_{eff} = -\frac{4G_F}{\sqrt{2}} V_{tb}^* V_{ts} \sum_{i=1}^{10} C_i(\mu) O_i(\mu), \quad (3)$$

where $O_i(\mu)$ ($i = 1, \dots, 10$) are the four-quark operators and $C_i(\mu)$ are the corresponding Wilson coefficients at the energy scale μ and the explicit expressions of these in the SM at NLO and NNLL are given in [47–58]. The operators responsible for $B_c \rightarrow D_s^* \ell^+ \ell^-$ are O_7 , O_9 and O_{10} and their form is given by

$$\begin{aligned} O_7 &= \frac{e^2}{16\pi^2} m_b (\bar{s} \sigma_{\mu\nu} P_R b) F^{\mu\nu}, \\ O_9 &= \frac{e^2}{16\pi^2} (\bar{s} \gamma_\mu P_L b) (\bar{\ell} \gamma^\mu \ell), \\ O_{10} &= \frac{e^2}{16\pi^2} (\bar{s} \gamma_\mu P_L b) (\bar{\ell} \gamma^\mu \gamma_5 \ell), \end{aligned} \quad (4)$$

with $P_{L,R} = (1 \pm \gamma_5)/2$.

The sequential fourth generation model with an additional up-type quark t' and down-type quark b' , a heavy charged lepton τ' and an associated neutrino ν' is a simple and non-supersymmetric extension of the SM, and as such does not add any new dynamics to the SM. Being a simple extension of the SM it retains all the properties of the SM where the new top quark t' like the other up-type quarks, contributes to $b \rightarrow s$ transition at the loop level. Therefore, the effect of fourth generation displays itself by changing the values of Wilson coefficients $C_7(\mu)$, $C_9(\mu)$ and C_{10} via the virtual exchange of fourth generation up-type quark t' which then take the form;

$$\lambda_t C_i \rightarrow \lambda_t C_i^{SM} + \lambda_{t'} C_i^{new}, \quad (5)$$

where $\lambda_f = V_{fb}^* V_{fs}$ and the explicit forms of the C_i 's can be obtained from the corresponding expressions of the Wilson coefficients in the SM by substituting $m_t \rightarrow m_{t'}$. By adding an extra family of quarks, the CKM matrix of the SM is extended by another row and column which now becomes 4×4 . The unitarity of which leads to

$$\lambda_u + \lambda_c + \lambda_t + \lambda_{t'} = 0.$$

Since $\lambda_u = V_{ub}^* V_{us}$ has a very small value compared to the others, therefore, we will ignore it. Then $\lambda_t \approx -\lambda_c - \lambda_{t'}$ and from Eq. (5) we have

$$\lambda_t C_i^{SM} + \lambda_{t'} C_i^{new} = -\lambda_c C_i^{SM} + \lambda_{t'} (C_i^{new} - C_i^{SM}). \quad (6)$$

One can clearly see that under $\lambda_{t'} \rightarrow 0$ or $m_{t'} \rightarrow m_t$ the term $\lambda_{t'} (C_i^{new} - C_i^{SM})$ vanishes which is the requirement of GIM mechanism. Taking the contribution of the t' quark in the loop the Wilson coefficients C_i 's can be written in the following form

$$\begin{aligned} C_7^{tot}(\mu) &= C_7^{SM}(\mu) + \frac{\lambda_{t'}}{\lambda_t} C_7^{new}(\mu), \\ C_9^{tot}(\mu) &= C_9^{SM}(\mu) + \frac{\lambda_{t'}}{\lambda_t} C_9^{new}(\mu), \\ C_{10}^{tot} &= C_{10}^{SM} + \frac{\lambda_{t'}}{\lambda_t} C_{10}^{new}, \end{aligned} \quad (7)$$

where we factored out $\lambda_t = V_{tb}^* V_{ts}$ term in the effective Hamiltonian given in Eq. (3) and the last term in these expressions corresponds to the contribution of the t' quark to the Wilson Coefficients. $\lambda_{t'}$ can be parameterized as:

$$\lambda_{t'} = |V_{t'b}^* V_{t's}| e^{i\phi_{sb}} \quad (8)$$

where ϕ_{sb} is the new CP odd phase.

In terms of the above Hamiltonian, the free quark decay amplitude for $b \rightarrow s \ell^+ \ell^-$ in SM4 can be derived as:

$$\begin{aligned} \mathcal{M}(b \rightarrow s \ell^+ \ell^-) &= -\frac{G_F \alpha}{\sqrt{2}\pi} V_{tb} V_{ts}^* \left\{ C_9^{tot}(\bar{s} \gamma_\mu P_L b)(\bar{\ell} \gamma^\mu \ell) + C_{10}^{tot}(\bar{s} \gamma_\mu P_L b)(\bar{\ell} \gamma^\mu \gamma_5 \ell) \right. \\ &\quad \left. - 2m_b C_7^{tot}(\bar{s} i \sigma_{\mu\nu} \frac{q^\nu}{q^2} P_R b)(\bar{\ell} \gamma^\mu \ell) \right\}, \end{aligned} \quad (9)$$

where q^2 is the square of momentum transfer. The operator O_{10} can not be induced by the insertion of four-quark operators because of the absence of the Z -boson in the effective theory. Therefore, the Wilson coefficient C_{10} does not renormalize under QCD corrections and hence it is independent on the energy scale. In addition to this, the above quark level decay amplitude can receive contributions from the matrix element of four-quark operators, $\sum_{i=1}^6 \langle \ell^+ \ell^- s | O_i | b \rangle$, which are usually absorbed into the effective Wilson coefficient $C_9^{SM}(\mu)$ and can usually be called C_9^{eff} , that one can decompose into the following three parts

$$C_9^{SM} = C_9^{eff}(\mu) = C_9(\mu) + Y_{SD}(z, s') + Y_{LD}(z, s'),$$

where the parameters z and s' are defined as $z = m_c/m_b$, $s' = q^2/m_b^2$. $Y_{SD}(z, s')$ describes the short-distance contributions from four-quark operators far away from the $c\bar{c}$ resonance regions, which can be calculated reliably in the perturbative theory. The long-distance contributions $Y_{LD}(z, s')$ from four-quark operators near the $c\bar{c}$ resonance cannot be calculated from first principles of QCD and are usually parameterized in the form of a phenomenological Breit-Wigner formula making use of the vacuum saturation approximation and quark-hadron duality. We will not incorporate the long-distance contributions in this work. The expressions for $Y_{SD}(z, s')$ can be manifestly written as [48]

$$\begin{aligned} Y_{SD}(z, s') &= h(z, s')(3C_1(\mu) + C_2(\mu) + 3C_3(\mu) + C_4(\mu) + 3C_5(\mu) + C_6(\mu)) \\ &\quad - \frac{1}{2}h(1, s')(4C_3(\mu) + 4C_4(\mu) + 3C_5(\mu) + C_6(\mu)) \\ &\quad - \frac{1}{2}h(0, s')(C_3(\mu) + 3C_4(\mu)) + \frac{2}{9}(3C_3(\mu) + C_4(\mu) + 3C_5(\mu) + C_6(\mu)), \end{aligned} \quad (10)$$

with

$$h(z, s') = -\frac{8}{9}\ln z + \frac{8}{27} + \frac{4}{9}x - \frac{2}{9}(2+x)|1-x|^{1/2} \begin{cases} \ln \left| \frac{\sqrt{1-x}+1}{\sqrt{1-x}-1} \right| - i\pi & \text{for } x \equiv 4z^2/s' < 1 \\ 2 \arctan \frac{1}{\sqrt{x-1}} & \text{for } x \equiv 4z^2/s' > 1 \end{cases},$$

$$h(0, s') = \frac{8}{27} - \frac{8}{9}\ln \frac{m_b}{\mu} - \frac{4}{9}\ln s' + \frac{4}{9}i\pi. \quad (11)$$

C. Parameterizations of matrix elements and form factors in QCD Sum Rules

The exclusive $B_c \rightarrow D_s^* \ell^+ \ell^-$ decay involves the hadronic matrix elements which can be obtained by sandwiching the quark level operators given in Eq. (9) between initial state B_c meson and final state D_s^* meson. These can be parameterized in terms of the form factors which are the scalar functions of the square of the four momentum transfer ($q^2 = (p - k)^2$). The non vanishing matrix elements for the process $B_c \rightarrow D_s^*$ can be parameterized in terms of the seven form factors as follows

$$\begin{aligned} \langle D_s^*(k, \varepsilon) | \bar{s} \gamma_\mu b | B_c(p) \rangle &= \frac{2A_V(q^2)}{M_{B_c} + M_{D_s^*}} \epsilon_{\mu\nu\alpha\beta} \varepsilon^{*\nu} p^\alpha k^\beta \\ \langle D_s^*(k, \varepsilon) | \bar{s} \gamma_\mu \gamma_5 b | B_c(p) \rangle &= i \left(M_{B_c^-} + M_{D_s^*} \right) \varepsilon^{*\mu} A_0(q^2) \\ &\quad - i \frac{A_+(q^2)}{M_{B_c} + M_{D_s^*}} (\varepsilon^* \cdot q) (p + k)^\mu \\ &\quad - i \frac{A_-(q^2)}{M_{B_c} + M_{D_s^*}} (\varepsilon^* \cdot q) q^\mu \end{aligned} \quad (12)$$

$$(13)$$

where p is the momentum of B_c meson and $\varepsilon(k)$ are the polarization vector (momentum) of the final state D_s^* meson.

In addition to the above form factors there are some penguin form factors, which we can write as

$$\begin{aligned} \langle D_s^*(k, \varepsilon) | \bar{s} \sigma_{\mu\nu} q^\nu b | B_c(p) \rangle &= 2iT_1(q^2) \epsilon_{\mu\nu\alpha\beta} \varepsilon^{*\nu} p^\alpha k^\beta \\ \langle D_s^*(k, \varepsilon) | \bar{s} \sigma_{\mu\nu} q^\nu \gamma_5 b | B_c(p) \rangle &= \left[\left(M_{B_c}^2 - M_{D_s^*}^2 \right) \varepsilon_\mu^* - (\varepsilon^* \cdot q)(p + k)_\mu \right] T_2(q^2) \\ &\quad + (\varepsilon^* \cdot q) \left[q_\mu - \frac{q^2}{M_{B_c}^2 - M_{D_s^*}^2} (p + k)_\mu \right] T_3(q^2). \end{aligned} \quad (14)$$

$$(15)$$

The form factors $A_V(q^2)$, $A_0(q^2)$, $A_+(q^2)$, $A_-(q^2)$, $T_1(q^2)$, $T_2(q^2)$, $T_3(q^2)$ are the non-perturbative quantities and to calculate them one has to rely on some non-perturbative approaches and in our numerical analysis we use the form factors calculated by using QCD Sum Rules [44]. The dependence of these form factors on square of the momentum transfer (q^2) can be written as

$$F(q^2) = \frac{F(0)}{1 + a \frac{q^2}{M_{B_c}^2} + b \frac{q^4}{M_{B_c}^4}}. \quad (16)$$

where the values of the parameters $F(0)$, α and β are given in Table II.

Now in terms of these form factors and from Eq. (9) it is straightforward to write the penguin amplitude as

$$\mathcal{M}^{\text{PENG}} = -\frac{G_F \alpha}{2\sqrt{2}\pi} V_{tb} V_{ts}^* [T_\mu^1 (\bar{l} \gamma^\mu l) + T_\mu^2 (\bar{l} \gamma^\mu \gamma_5 l)]$$

TABLE II: $B_c \rightarrow D_s^*$ form factors corresponding to penguin contributions in the QCD Sum Rules. $F(0)$ denotes the value of form factors at $q^2 = 0$ while a and b are the parameters in the parameterizations shown in Eq. (16)[44].

$F(q^2)$	$F(0)$	a	b
$A_V(q^2)$	0.54 ± 0.018	-1.28	-0.23
$A_0(q^2)$	0.30 ± 0.017	-0.13	-0.18
$A_+(q^2)$	0.36 ± 0.013	-0.67	-0.066
$A_-(q^2)$	-0.57 ± 0.04	-1.11	-0.14
$T_1(q^2)$	0.31 ± 0.017	-1.28	-0.23
$T_2(q^2)$	0.33 ± 0.016	-0.10	-0.097
$T_3(q^2)$	0.29 ± 0.034	-0.91	0.007

where

$$T_\mu^1 = f_1(q^2)\epsilon_{\mu\nu\alpha\beta}\varepsilon^{*\nu}p^\alpha k^\beta - if_2(q^2)\varepsilon_\mu^* + if_3(q^2)(\varepsilon^* \cdot q)P_\mu \quad (17)$$

$$T_\mu^2 = f_4(q^2)\epsilon_{\mu\nu\alpha\beta}\varepsilon^{*\nu}p^\alpha k^\beta - if_5(q^2)\varepsilon_\mu^* + if_6(q^2)(\varepsilon^* \cdot q)P_\mu + if_0(q^2)(\varepsilon^* \cdot q)q_\mu \quad (18)$$

The functions f_0 to f_6 in Eq. (17) and Eq. (18) are known as auxiliary functions, which contain both long distance (form factors) and short distance (Wilson coefficients) effects and these can be written as

$$\begin{aligned}
f_1(q^2) &= 4(m_b + m_s)\frac{C_7^{tot}}{q^2}T_1(q^2) + 2C_9^{tot}\frac{A_V(q^2)}{M_{B_c} + M_{D_s^*}} \\
f_2(q^2) &= \frac{2C_7^{tot}}{q^2}(m_b - m_s)T_2(q^2)\left(M_{B_c}^2 - M_{D_s^*}^2\right) + C_9^{tot}A_0(q^2)(M_{B_c} + M_{D_s^*}) \\
f_3(q^2) &= \left[4\frac{C_7^{tot}}{q^2}(m_b - m_s)\left(T_2(q^2) + q^2\frac{T_3(q^2)}{(M_{B_c}^2 - M_{D_s^*}^2)}\right) + C_9^{tot}\frac{A_+(q^2)}{M_{B_c} + M_{D_s^*}}\right] \\
f_4(q^2) &= C_{10}^{tot}\frac{2A_V(q^2)}{M_{B_c} + M_{D_s^*}} \\
f_5(q^2) &= C_{10}^{tot}A_0(q^2)(M_{B_c} + M_{D_s^*}) \\
f_6(q^2) &= C_{10}^{tot}\frac{A_+(q^2)}{M_{B_c} + M_{D_s^*}} \\
f_0(q^2) &= C_{10}^{tot}\frac{A_-(q^2)}{M_{B_c} + M_{D_s^*}} \quad (19)
\end{aligned}$$

III. PHYSICAL OBSERVABLES FOR $B_c \rightarrow D_s^*\ell^+\ell^-$

In this section we will present the calculations of the physical observables like the decay rates, leptons forward-backward asymmetry, the helicity fractions of D_s^* meson and the final state lepton polarizations. We use both the weak annihilation (WA) amplitude and the penguin amplitude to study these observables.

A. The Differential Decay Rate of $B_c \rightarrow D_s^* \ell^+ \ell^-$

In the rest frame of B_c meson the differential decay width of $B_c \rightarrow D_s^* \ell^+ \ell^-$ can be written as

$$\frac{d\Gamma(B_c \rightarrow D_s^* \ell^+ \ell^-)}{dq^2} = \frac{1}{(2\pi)^3} \frac{1}{32M_{B_c}^3} \int_{-u(q^2)}^{+u(q^2)} du |\mathcal{M}|^2 \quad (20)$$

where

$$\begin{aligned} \mathcal{M} &= \mathcal{M}^{\text{WA}} + \mathcal{M}^{\text{PENG}} \\ q^2 &= (p_{l^+} + p_{l^-})^2 \\ u &= (p - p_{l^-})^2 - (p - p_{l^+})^2 \end{aligned} \quad (21)$$

Now the limits on q^2 and u are

$$4m_l^2 \leq q^2 \leq (M_{B_c} - M_{D_s^*})^2 \quad (22)$$

$$-u(q^2) \leq u \leq u(q^2) \quad (23)$$

with

$$u(q^2) = \sqrt{\lambda \left(1 - \frac{4m_l^2}{q^2} \right)} \quad (24)$$

and

$$\lambda \equiv \lambda(M_{B_c}^2, M_{D_s^*}^2, q^2) = M_{B_c}^4 + M_{D_s^*}^4 + q^4 - 2M_{B_c}^2 M_{D_s^*}^2 - 2M_{D_s^*}^2 q^2 - 2q^2 M_{B_c}^2$$

Here m_l corresponds to the mass of the lepton which for our case are the μ and τ . The total decay rate of $B_c \rightarrow D_s^* \ell^+ \ell^-$ can be expressed in terms of WA, penguin amplitude and their interference, which takes the form [41]

$$\frac{d\Gamma}{dq^2} = \frac{d\Gamma^{\text{WA}}}{dq^2} + \frac{d\Gamma^{\text{PENG}}}{dq^2} + \frac{d\Gamma^{\text{WA-PENG}}}{dq^2} \quad (25)$$

with

$$\frac{d\Gamma^{\text{WA}}}{dq^2} = \frac{G_F^2 |V_{cb} V_{cs}^*|^2 \alpha^2}{2^{11} \pi^5 3 M_{B_c}^3 M_{D_s^*}^2 q^2} u(q^2) \times g(q^2) \quad (26)$$

$$\frac{d\Gamma^{\text{PENG}}}{dq^2} = \frac{G_F^2 |V_{tb} V_{ts}^*|^2 \alpha^2}{2^{11} \pi^5 3 M_{B_c}^3 M_{D_s^*}^2 q^2} u(q^2) \times h(q^2) \quad (27)$$

$$\frac{d\Gamma^{\text{WA-PENG}}}{dq^2} = \frac{G_F^2 |V_{cb} V_{cs}^*| |V_{tb} V_{ts}^*| \alpha^2}{2^{11} \pi^5 3 M_{B_c}^3 M_{D_s^*}^2 q^2} u(q^2) \times I(q^2). \quad (28)$$

The function $u(q^2)$ is defined in Eq. (24) and $g(q^2)$, $h(q^2)$ and $I(q^2)$ are

$$\begin{aligned} g(q^2) &= \frac{1}{2} (2m_l^2 + q^2) \kappa^2 \left[8\lambda M_{D_s^*}^2 q^2 \left(F_V^{D_s^*}(q^2) \right)^2 + \left(F_A^{D_s^*}(q^2) \right)^2 [12M_{D_s^*}^2 q^2 (\lambda + 4M_{B_c}^2 q^2) + \lambda^2 + \lambda(\lambda + 4q^2 M_{D_s^*}^2 + 4q^4)] \right] \\ h(q^2) &= 24 |f_0(q^2)|^2 m_l^2 M_{D_s^*}^2 \lambda + 8M_{D_s^*}^2 q^2 \lambda (2m_l^2 + q^2) |f_1(q^2)|^2 - (4m_l^2 - q^2) |f_4(q^2)|^2 \\ &\quad + \lambda (2m_l^2 + q^2) \left| f_2(q^2) + (M_{B_c}^2 - M_{D_s^*}^2 - q^2) f_3(q^2) \right|^2 - (4m_l^2 - q^2) \left| f_5(q^2) + (M_{B_c}^2 - M_{D_s^*}^2 - q^2) f_6(q^2) \right|^2 \\ &\quad + 4M_{D_s^*}^2 q^2 [(2m_l^2 + q^2)(3 |f_2(q^2)|^2 - \lambda |f_3(q^2)|^2) - (4m_l^2 - q^2)(3 |f_5(q^2)|^2 - \lambda |f_6(q^2)|^2)] \\ I(q^2) &= 2\kappa [f_2(q^2) F_A^{D_s^*}(q^2) q^2 (2m_l^2 + q^2) (\lambda + 6M_{D_s^*}^2 (M_{B_c}^2 - M_{D_s^*}^2 + q^2)) \\ &\quad - (\lambda (2f_1(q^2) F_V^{D_s^*}(q^2) M_{D_s^*}^2 q^4 + f_3(q^2) F_A^{D_s^*}(q^2) (2m_l^2 + q^2) (\lambda + q^4 + 4M_{B_c} M_{D_s^*}))]. \end{aligned} \quad (29)$$

where

$$\kappa = \frac{8\pi^2 M_{D_s^*} f_{B_c} f_{D_s^*} |V_{tb} V_{ts}^*|}{(m_c^2 - m_s^2) q^2 |V_{cb} V_{cs}^*|}. \quad (30)$$

B. Forward-Backward Asymmetry of $B \rightarrow D_s^* \ell^+ \ell^-$ decay

In this section we investigate the forward-backward asymmetry (FBA) of leptons. The measurement of the FBA at LHC is significant due to the minimal dependence on the form factors [49], and as such this observable is of greater importance to check the more clear signals of any NP than the other observables such as branching ratio etc. In the context of fourth generation, the FBA can also play a crucial role because it is driven by the loop top quark and so it is sensitive to the fourth generation up type quark t' [50]. Now to calculate the forward-backward asymmetry, we consider the following double differential decay rate formula for the process $B_c \rightarrow D_s^* \ell^+ \ell^-$

$$\frac{d^2\Gamma(q^2, \cos\theta)}{dq^2 d\cos\theta} = \frac{1}{(2\pi)^3} \frac{1}{32m_B^3} u(q^2) |\mathcal{M}_{B_c \rightarrow D_s^* \ell^+ \ell^-}|^2, \quad (31)$$

where θ is the angle between the momentum of B_c meson and ℓ^- in the dilepton rest frame. Following Refs. [50], the differential and normalized FBAs for the semi-leptonic decay $B_c \rightarrow D_s^* \ell^+ \ell^-$ are defined as

$$\frac{dA_{FB}(q^2)}{dq^2} = \int_0^1 d\cos\theta \frac{d^2\Gamma(q^2, \cos\theta)}{dq^2 d\cos\theta} - \int_{-1}^0 d\cos\theta \frac{d^2\Gamma(q^2, \cos\theta)}{dq^2 d\cos\theta} \quad (32)$$

and

$$A_{FB}(q^2) = \frac{\int_0^1 d\cos\theta \frac{d^2\Gamma(q^2, \cos\theta)}{dq^2 d\cos\theta} - \int_{-1}^0 d\cos\theta \frac{d^2\Gamma(q^2, \cos\theta)}{dq^2 d\cos\theta}}{\int_0^1 d\cos\theta \frac{d^2\Gamma(q^2, \cos\theta)}{dq^2 d\cos\theta} + \int_{-1}^0 d\cos\theta \frac{d^2\Gamma(q^2, \cos\theta)}{dq^2 d\cos\theta}}. \quad (33)$$

Following the procedure used for the differential decay rate, one can easily get the expression for the forward-backward asymmetry which can be written as

$$\mathcal{A}_{FB} = \frac{1}{d\Gamma/dq^2} \frac{G_F^2 \alpha^2}{2^{11} \pi^5 m_B^3} |V_{tb} V_{ts}^*|^2 q^2 u(q^2) [2\text{Re}[f_4] \kappa F_A^{D_s^*}(q^2) (M_{B_c}^2 - M_{D_s^*}^2 + q^2) + 4\text{Re}[f_2^* f_4 + f_1^* f_5] + 2\text{Re}[f_5] \kappa F_V^{D_s^*}(q^2)] \quad (34)$$

where κ is defined in Eq. (30) and $d\Gamma/dq^2$ is given in Eq. (25).

C. Helicity Fractions of D_s^* in $B_c \rightarrow D_s^* \ell^+ \ell^-$ decay

We now discuss helicity fractions of D_s^* meson in $B_c \rightarrow D_s^* \ell^+ \ell^-$ decay which are almost independent of the uncertainties arising due to form factors and other input parameters. Therefore, the study of these observables will provide us a good testing ground for the SM4. The explicit expression of the decay rate for $B_c^- \rightarrow D_s^* \ell^+ \ell^-$ decay in terms of longitudinal (Γ_L) and transverse (Γ_T) components of decay rate can be written as [41]

$$\frac{d\Gamma_L(q^2)}{dq^2} = \frac{d\Gamma_L^{\text{WA}}(q^2)}{dq^2} + \frac{d\Gamma_L^{\text{PENG}}(q^2)}{dq^2} + \frac{d\Gamma_L^{\text{WA-PENG}}(q^2)}{dq^2} \quad (35)$$

$$\frac{d\Gamma_{\pm}(q^2)}{dq^2} = \frac{d\Gamma_{\pm}^{\text{WA}}(q^2)}{dq^2} + \frac{d\Gamma_{\pm}^{\text{PENG}}(q^2)}{dq^2} + \frac{d\Gamma_{\pm}^{\text{WA-PENG}}(q^2)}{dq^2} \quad (36)$$

$$\frac{d\Gamma_T(q^2)}{dq^2} = \frac{d\Gamma_+(q^2)}{dq^2} + \frac{d\Gamma_-(q^2)}{dq^2}. \quad (37)$$

where

$$\frac{d\Gamma_L^{\text{WA}}(q^2)}{dq^2} = \frac{G_F^2 |V_{cb} V_{cs}^*|^2 \alpha^2 u(q^2)}{2^{11} \pi^5 M_{B_c}^3} \times \frac{1}{3} A_L^{\text{WA}} \quad (38)$$

$$\frac{d\Gamma_L^{\text{PENG}}(q^2)}{dq^2} = \frac{G_F^2 |V_{tb} V_{ts}^*|^2 \alpha^2 u(q^2)}{2^{11} \pi^5 M_{B_c}^3} \times \frac{1}{3} A_L^{\text{PENG}} \quad (39)$$

$$\frac{d\Gamma_L^{\text{WA-PENG}}(q^2)}{dq^2} = \frac{G_F^2 |V_{cb} V_{cs}^*| |V_{tb} V_{ts}^*| \alpha^2 u(q^2)}{2^{11} \pi^5 M_{B_c}^3} \times \frac{1}{3} A_L^{\text{WA-PENG}} \quad (40)$$

$$\frac{d\Gamma_{\pm}^{\text{WA}}(q^2)}{dq^2} = \frac{G_F^2 |V_{cb} V_{cs}^*|^2 \alpha^2 u(q^2)}{2^{11} \pi^5 M_{B_c}^3} \times \frac{2}{3} A_{\pm}^{\text{WA}} \quad (41)$$

$$\quad (42)$$

$$\frac{d\Gamma_{\pm}^{\text{PENG}}(q^2)}{dq^2} = \frac{G_F^2 |V_{tb} V_{ts}^*|^2 \alpha^2 u(q^2)}{2^{11} \pi^5 M_{B_c}^3} \times \frac{4}{3} A_{\pm}^{\text{PENG}} \quad (43)$$

$$\frac{d\Gamma_{\pm}^{\text{WA-PENG}}(q^2)}{dq^2} = \frac{G_F^2 |V_{cb} V_{cs}^*| |V_{tb} V_{ts}^*| \alpha^2 u(q^2)}{2^{11} \pi^5 M_{B_c}^3} \times \frac{2}{3} A_{\pm}^{\text{WA-EP}}. \quad (44)$$

The different functions appearing in above equation can be expressed in terms of auxiliary functions (c.f. Eq. (19))

as

$$\begin{aligned} A_L^{\text{WA}} &= \frac{\kappa^2}{4q^2 M_{D_s^*}^2} \left[\left(F_V^{D_s^*}(q^2) \right)^2 \left\{ q^2 \lambda (\lambda + 4q^2 M_{D_s^*}^2) - 4M^2 \lambda (2\lambda + 8q^2 M_{D_s^*}^2) - q^2 (M_{B_c}^2 - M_{D_s^*}^2 - q^2)^2 (\lambda - 2u^2(q^2)) \right\} \right. \\ &\quad + \left(F_A^{D_s^*}(q^2) \right)^2 \left\{ 12\lambda q^2 ((M_{B_c}^2 - M_{D_s^*}^2)^2 - M_{D_s^*}^2) - \lambda^2 (q^2 - 4m^2) + q^2 (8q^2 M_{D_s^*}^2 - \lambda) (M_{B_c}^2 - M_{D_s^*}^2 + q^2)^2 \right. \\ &\quad \left. \left. - 2u^2(q^2) q^2 ((M_{B_c}^2 - M_{D_s^*}^2)^2 + q^4) + 4m^2 ((M_{B_c}^2 - M_{D_s^*}^2)^2 - q^4)^2 \right\} \right] \\ A_L^{\text{PENG}} &= \frac{1}{2M_{D_s^*}^2 q^2} [24 |f_0(q^2)|^2 m_l^2 M_{D_s^*}^2 \lambda + (2m_l^2 + q^2) \left| (M_{B_c}^2 - M_{D_s^*}^2 - q^2) f_2(q^2) + \lambda f_3(q^2) \right|^2 \\ &\quad + (q^2 - 4m_l^2) \left| (M_{B_c}^2 - M_{D_s^*}^2 - q^2) f_5(q^2) + \lambda f_6(q^2) \right|^2] \\ A_L^{\text{WA-PENG}} &= \frac{\kappa}{q^2 M_{D_s^*}^2} \left[Re(f_1(q^2) F_V^{D_s^*}(q^2)) \left\{ (\lambda + 4M_{D_s^*}^2 q^2) (8m^2 \sqrt{\lambda} + q^2 (2u(q^2) - \sqrt{\lambda})) - 4M_{D_s^*}^2 q^2 \lambda \right\} \right. \\ &\quad + Re(f_2(q^2) F_A^{D_s^*}(q^2)) \left\{ q^2 u^2(q^2) (M_{B_c}^2 - M_{D_s^*}^2 - q^2) + 6q^2 \lambda (M_{D_s^*}^2 - M_{B_c}^2) \right. \\ &\quad \left. \left. + q^2 (\lambda - 8q^2 M_{D_s^*}^2) (M_{B_c}^2 - M_{D_s^*}^2 + q^2) - 4m^2 q^2 (4q^2 M_{D_s^*}^2 + \lambda) \right\} \right. \\ &\quad \left. + Re(f_3(q^2) F_A^{D_s^*}(q^2)) \left\{ \lambda^2 (4m^2 - q^2) + q^4 (q^2 u(q^2) \sqrt{\lambda} - 6\lambda (M_{B_c}^2 + M_{D_s^*}^2)) + q^2 (M_{B_c}^2 - M_{D_s^*}^2) (6\lambda - u^2(q^2)) \right\} \right] \\ A_{\pm}^{\text{WA}} &= \kappa^2 \left[(2m^2 + q^2) \left[\lambda \left(F_V^{D_s^*}(q^2) \right)^2 + \left(F_A^{D_s^*}(q^2) \right)^2 (\lambda + 4M_{D_s^*}^2 q^2) \right] \right] \\ A_{\pm}^{\text{PENG}} &= (q^2 - 4m_l^2) \left| f_5(q^2) \mp \sqrt{\lambda} f_4(q^2) \right|^2 + (q^2 + 2m_l^2) \left| f_2(q^2) \pm \sqrt{\lambda} f_1(q^2) \right|^2 \\ A_{\pm}^{\text{WA-PENG}} &= -\kappa \left\{ 2\sqrt{\lambda} (q^2 - 4m^2) Re(f_2(q^2) F_V^{D_s^*}(q^2)) + 4\lambda (q^2 + 2m^2) Re(f_1(q^2) F_V^{D_s^*}(q^2)) \right. \\ &\quad \left. \pm 2(q^2 + 2m^2) (M_{B_c}^2 - M_{D_s^*}^2 + q^2) [2Re(f_1(q^2) F_A^{D_s^*}(q^2))] \sqrt{\lambda} \mp 2Re(f_2(q^2) F_V^{D_s^*}(q^2)) \right\} \end{aligned} \quad (45)$$

Finally the longitudinal and transverse helicity amplitude becomes

$$\begin{aligned} f_L(q^2) &= \frac{d\Gamma_L(q^2)/dq^2}{d\Gamma(q^2)/dq^2} \\ f_{\pm}(q^2) &= \frac{d\Gamma_{\pm}(q^2)/dq^2}{d\Gamma(q^2)/dq^2} \\ f_T(q^2) &= f_+(q^2) + f_-(q^2) \end{aligned} \quad (46)$$

so that the sum of the longitudinal and transverse helicity amplitudes is equal to one i.e. $f_L(q^2) + f_T(q^2) = 1$ for each value of q^2 .

D. Lepton Polarization asymmetries of $B_c \rightarrow D_s^* \ell^+ \ell^-$

In the rest frame of the lepton ℓ^- , the unit vectors along longitudinal, normal and transversal component of the ℓ^- can be defined as:

$$\begin{aligned} s_L^{-\mu} &= (0, \vec{e}_L) = \left(0, \frac{\vec{p}_-}{|\vec{p}_-|}\right), \\ s_N^{-\mu} &= (0, \vec{e}_N) = \left(0, \frac{\vec{k} \times \vec{p}_-}{|\vec{k} \times \vec{p}_-|}\right), \\ s_T^{-\mu} &= (0, \vec{e}_T) = (0, \vec{e}_N \times \vec{e}_L), \end{aligned} \quad (47)$$

where \vec{p}_- and \vec{k} are the respective three-momenta of the lepton ℓ^- and D_s^* meson in the center mass (CM) frame of $l^+ l^-$ system. Lorentz transformation is used to boost the longitudinal component of the lepton polarization to the CM frame of the lepton pair as

$$(s_L^{-\mu})_{CM} = \left(\frac{|\vec{p}_-|}{m_l}, \frac{E_l \vec{p}_-}{m_l |\vec{p}_-|}\right) \quad (48)$$

where E_l and m_l are the energy and mass of the lepton. The normal and transverse components remain unchanged under the Lorentz boost.

The longitudinal (P_L), normal (P_N) and transverse (P_T) polarizations of lepton can be defined as:

$$P_i^{(\mp)}(q^2) = \frac{\frac{d\Gamma}{dq^2}(\vec{\xi}^{\mp} = \vec{e}^{\mp}) - \frac{d\Gamma}{dq^2}(\vec{\xi}^{\mp} = -\vec{e}^{\mp})}{\frac{d\Gamma}{dq^2}(\vec{\xi}^{\mp} = \vec{e}^{\mp}) + \frac{d\Gamma}{dq^2}(\vec{\xi}^{\mp} = -\vec{e}^{\mp})} \quad (49)$$

where $i = L, N, T$ and $\vec{\xi}^{\mp}$ is the spin direction along the leptons l^{\mp} . The differential decay rate for polarized lepton l^{\mp} in $B_c \rightarrow D_s^* \ell^+ \ell^-$ decay along any spin direction $\vec{\xi}^{\mp}$ is related to the unpolarized decay rate (25) with the following relation

$$\frac{d\Gamma(\vec{\xi}^{\mp})}{dq^2} = \frac{1}{2} \left(\frac{d\Gamma}{dq^2} \right) [1 + (P_L^{\mp} \vec{e}_L^{\mp} + P_N^{\mp} \vec{e}_N^{\mp} + P_T^{\mp} \vec{e}_T^{\mp}) \cdot \vec{\xi}^{\mp}]. \quad (50)$$

Using these inputs we can achieve the expressions of longitudinal, normal and transverse lepton polarizations for $B_c \rightarrow D_s^* \ell^+ \ell^-$ decays. The expression of the numerator of longitudinal lepton polarization is

$$\begin{aligned} P_L(q^2) \propto & \frac{4\lambda}{3M_{D_s^*}^2} \sqrt{\frac{q^2 - 4m_l^2}{q^2}} \times \left\{ 2\text{Re}(f_2 f_5^*) + \lambda \text{Re}(f_3 f_6^*) + 4\sqrt{q^2} \text{Re}(f_1 f_4^*) \left(1 + \frac{12q^2 M_{D_s^*}^2}{\lambda} \right) \right. \\ & \left. + (-M_{B_c}^2 + M_{D_s^*}^2 + q^2) [\text{Re}(f_3 f_5^*) + \text{Re}(f_2 f_6^*)] \right\}. \end{aligned} \quad (51)$$

Similarly the numerator of normal lepton polarization can be written as

$$P_N(q^2) \propto \frac{m_l \pi}{M_{D_s^*}^2} \sqrt{\frac{\lambda}{q^2}} \times \left\{ -\lambda q^2 \text{Re}(f_3 f_0^*) + \lambda (M_{B_c}^2 - M_{D_s^*}^2) \text{Re}(f_3 f_6^*) - \lambda \text{Re}(f_3 f_5^*) \right. \\ \left. + (-M_{B_c}^2 + M_{D_s^*}^2 + q^2) [q^2 \text{Re}(f_2 f_0^*)] \right. \\ \left. - 8q^2 M_{D_s^*}^2 \text{Re}(f_1 f_2^*) + \kappa^2 F_V^{D_s^*} F_A^{D_s^*} (M_{B_c}^2 - M_{D_s^*}^2 + q^2) \right\} \quad (52)$$

and that of the transverse leptons polarization is given by

$$P_T(q^2) \propto i \frac{m_l \pi \sqrt{\left(q^2 - \frac{4m_l^2}{q^2}\right)} \lambda}{M_{D_s^*}^2} \left\{ M_{B_c}^2 \text{Im}(f_5 f_6^*) + M_{D_s^*} [4\text{Im}(f_2 f_4^*) + 4\text{Im}(f_1 f_5^*) + 3\text{Im}(f_5 f_6^*)] \right. \\ \left. + (-M_{B_c}^2 + M_{D_s^*}^2 + q^2) [\text{Im}(f_0 f_5^*)] + \text{Im}(f_5 f_6^*) + \text{Im}(f_4 f_6^*) + 2\kappa \text{Im}[f_5] F_V^{D_s^*} + 2\kappa M_{D_s^*} \text{Im}[f_4] F_A^{D_s^*} \right\} \quad (53)$$

where κ is defined in Eq. (30) along with auxiliary functions f_0, \dots, f_6 and the form factors $F_{V,A}^{D_s^*}$ are the ones defined in Eq. (19) and Eq. (2), respectively. Here we have dropped out the constant factors which are understood.

IV. NUMERICAL ANALYSIS:

In this section, we would like to present the numerical analysis of decay rates, FBAs of leptons, helicity fractions of final state D_s^* meson and different lepton polarization asymmetries both in the SM and SM4. The numerical values of Wilson coefficients and other input parameters used in our analysis are collected in Tables II and III. It has already

TABLE III: Values of input parameters used in our numerical analysis

$G_F = 1.166 \times 10^{-2} \text{ GeV}^{-2}$	$ V_{ts} = 41.61_{-0.80}^{+0.10} \times 10^{-3}$
$ V_{tb} = 0.9991$	$m_b = (4.68 \pm 0.03) \text{ GeV}$
$m_c(m_c) = 1.275_{-0.015}^{+0.015} \text{ GeV}$	$m_s(1 \text{ GeV}) = (142 \pm 28) \text{ MeV}$
$M_{B_c} = 6.26 \text{ GeV}$	$M_{D_s^*} = 2.12 \text{ GeV}$
$f_{B_c} = 0.35 \text{ GeV}$	$f_{D_s^*} = 0.30 \text{ GeV}$

TABLE IV: The Wilson coefficients $C_i(\mu)$ at the scale $\mu \sim m_b$ in the SM.

C_1	C_2	C_3	C_4	C_5	C_6	C_7	C_9	C_{10}
1.107	-0.248	-0.011	-0.026	-0.007	-0.031	-0.313	4.344	-4.669

been mentioned that in B_c to D_s^* transition the WA contributions are proportional to $V_{cb}V_{cs}^*$ and hence can not be ignored like in the ordinary $B_{u,d,s}$ to light meson decays. Using the values of the form factors given in Table I and II along with the value of input parameters from Tables III and IV, the numerical result of the branching ratios for the decays $B_c \rightarrow D_s^* \ell^+ \ell^-$ containing contributions from penguin, WA and both amplitudes are given in Table V [41]. Here, one can see that the WA contribution to the branching ratio for $B_c \rightarrow D_s^* \mu^+ \mu^-$ decay is almost an order of magnitude larger than the penguin ones therefore, we have to include it in the analysis of $B_c \rightarrow D_s^* \ell^+ \ell^-$ in SM4.

TABLE V: Branching ratio for $B_c \rightarrow D_s^* \mu^+ \mu^- (\tau^+ \tau^-)$ decay using form factors calculated in QCD sum rules [43, 44].

$BR^{(\text{PENG})}(B_c \rightarrow D_s^* \mu^+ \mu^- (\tau^+ \tau^-))$	$BR^{(\text{WA})}(B_c \rightarrow D_s^* \mu^+ \mu^- (\tau^+ \tau^-))$	$BR^{(\text{Total})}(B_c \rightarrow D_s^* \mu^+ \mu^- (\tau^+ \tau^-))$
$2.57 \times 10^{-7} (1.13 \times 10^{-8})$	$2.20 \times 10^{-6} (0.35 \times 10^{-9})$	$2.46 \times 10^{-6} (1.49 \times 10^{-8})$

Regarding the parameters of the SM4, recently CDF collaboration has given the lower bound on the mass of the t' quark to be $m_{t'} \geq 335$ GeV at 95% CL [59]. These bounds are little higher than the ones quoted in Ref. [60] of $m_{t'} \gtrsim 256$ GeV. On the other hand, the perturbativity of the Yukawa coupling implies that $m_{t'} \lesssim \sqrt{2\pi} \langle v \rangle \approx 600$ GeV, where $\langle v \rangle$ is the vacuum expectation value of the Higgs boson [61]. Thus, the mass $m_{t'}$ is constrained in a band, $m_{t'} = 335 - 600$ GeV, which increases the predictability of SM4. Keeping in the view that these bounds will be considerably improved at LHC, we set $m_{t'} = 300 - 600$ GeV in our numerical calculation. In addition to the masses of the sequential fourth generation of quarks the other important parameters are the CKM4 matrix elements, where $|V_{t's}|$ and $|V_{t'b}|$ are of the main interest for present study. The experimental upper bounds on these CKM matrix elements are $|V_{t's}| < 0.11$ and $|V_{t'b}| < 0.12$ [62, 63]. By taking the CKM unitarity condition, $\sum_i V_{is}^* V_{ib}$, ($i = u, c, t, t'$) together with the present measurements of the 3×3 CKM matrix [64], the bounds for CKM4 matrix elements are obtained to be [63, 65]

$$|V_{t's}^* V_{t'b}| \leq 1.5 \times 10^{-2}. \quad (54)$$

Incorporating these constraints on the fourth generation parameter space, the NP effects origination from SM4 on different physical observables are shown in Figs. 1-14.

Figs. 1 and 2 depict the variation of the differential branching ratio of $B_c \rightarrow D_s^* \mu^+ \mu^- (\tau^+ \tau^-)$ decays with q^2 both in the SM and in the sequential fourth generation model (SM4). These figures indicate that the values of the differential branching ratios are enhanced sizably with increase in the values of fourth generation parameters $m_{t'}$ and $|V_{t'b}^* V_{t's}|$. These new physics effects are prominent in whole q^2 region both for the μ and τ as the final state leptons which is due to the fact that at small value of q^2 the dominant contribution comes from $C_7^{tot}(\mu)$ whereas for the large value of q^2 the major contribution is from the Z exchange i.e., C_{10}^{tot} , which is sensitive to the mass of the fourth generation quark ($m_{t'}$).

As an exclusive decay, the new physics effects in the branching ratios are usually masked by the uncertainties involved in different input parameters where form factors are the major contributors. However, for the present case the new physics effects are very prominent and lies well separated from the SM values even in the error bounds. However there exists some other observables which have very mild dependence on the choice of the form factors. Among them the zero position of the forward-backward asymmetry, helicity fractions of final state D_s^* meson and the different lepton polarization asymmetries, are almost free from the hadronic uncertainties, in particular at low q^2 region, and hence serve as an important tool to study NP.

Fig. 3 describes the behavior of the forward-backward asymmetry of $B_c \rightarrow D_s^* \mu^+ \mu^-$ with q^2 . Here one can see that the value of the forward-backward asymmetry passes from the zero at a particular value of q^2 both in the SM as well as in the SM4. This is because of the destructive interference between the photon penguin (C_7^{eff}) and the Z penguin (C_9^{eff}). It has already been mentioned that the SM4 effects display themselves in the Wilson coefficients, therefore, one expects that both the zero crossing of FBA and its magnitude will be different from SM. This fact is illustrated in 3. We can see that the value of the forward-backward asymmetry decreases from the SM value but the position of

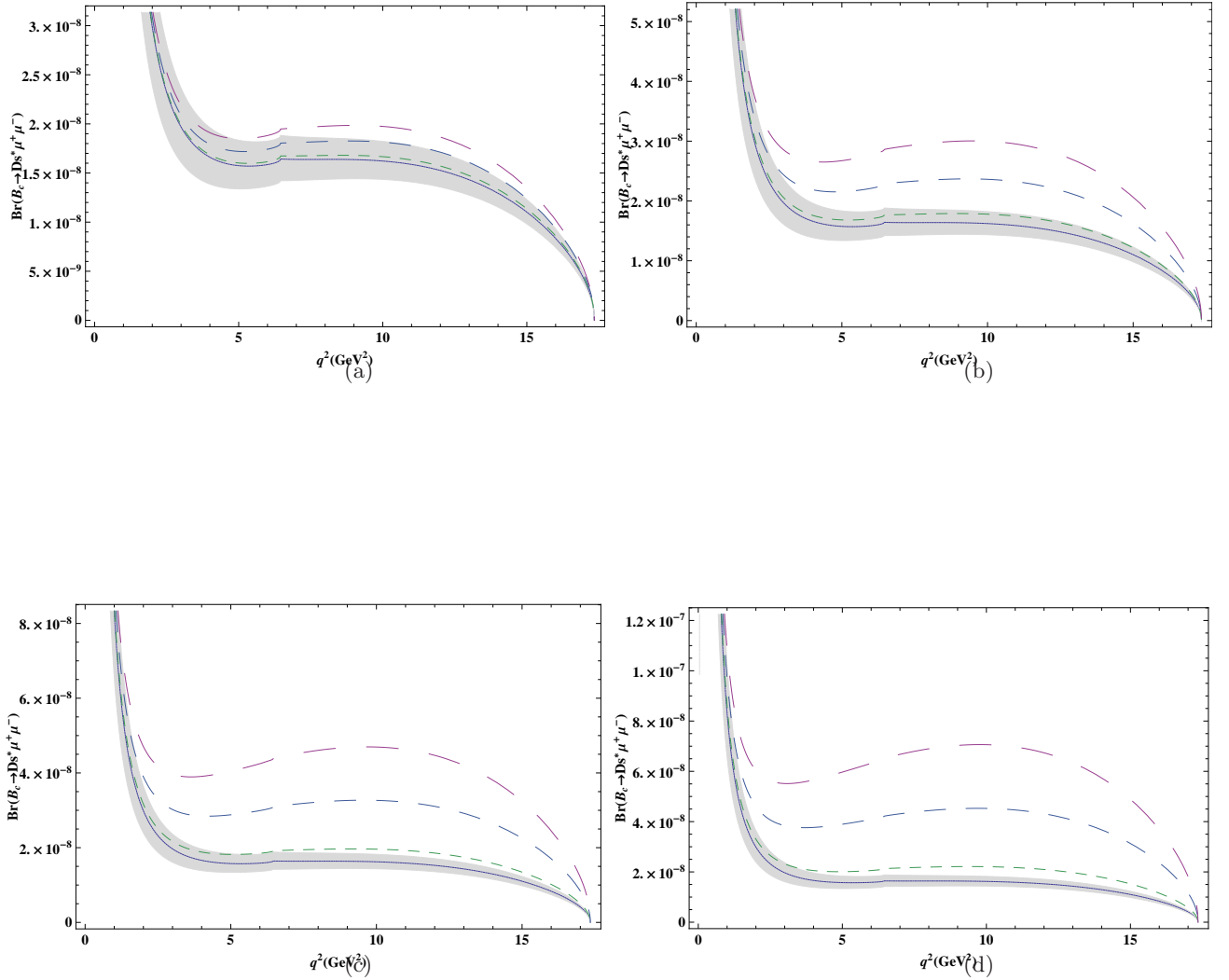


FIG. 1: The dependence of branching ratio of $B_c \rightarrow D_s^* \mu^+ \mu^-$ on q^2 for different values of $m_{t'}$ and $|V_{t'b}^* V_{t's}|$. $|V_{t'b}^* V_{t's}| = 0.003, 0.006, 0.009$ and 0.012 in (a), (b), (c) and (d) respectively. In all the graphs, the solid line corresponds to the SM, dashed, medium dashed and long dashed lines are for $m_{t'} = 300$ GeV, 500 GeV and 600 GeV respectively.

zero crossing remains the same for the low value of SM4 parameters ($m_{t'}, |V_{t'b}^* V_{t's}|$) (c.f. Fig. 3(a,b)). However, at the large value of the CKM4 matrix elements and the mass $m_{t'}$ the zero position is shifted to the left which makes it an important candidate for the search of SM4 effects.

Now for the $B_c \rightarrow D_s^* \tau^+ \tau^-$ decay the FBA is presented in Fig. 4. In this case the crossing of the FBA is absent both in the SM and in the SM4, however, there is a significant deviation in its magnitude for large values of the SM4 parameters. As the magnitude of the forward-backward asymmetry is also an important tool to establish the NP, therefore, the experimental study of it will give us some distinguishing effects of SM4.

Another handy tool to explore the NP is the study of helicity fractions of the final state meson which are associated with its spin. In literature, there exist some studies on the helicity fractions for the case of vector (K^*, D_s^*) and $K_1(1270, 1400)$ mesons both in the SM and in some NP scenarios, whereas for the K^* there exists some experimental observations too [41, 66]. It is therefore legitimate to study SM4 contributions to the helicity fractions of D_s^* meson in $B_c \rightarrow D_s^* \ell^+ \ell^-$ decays.

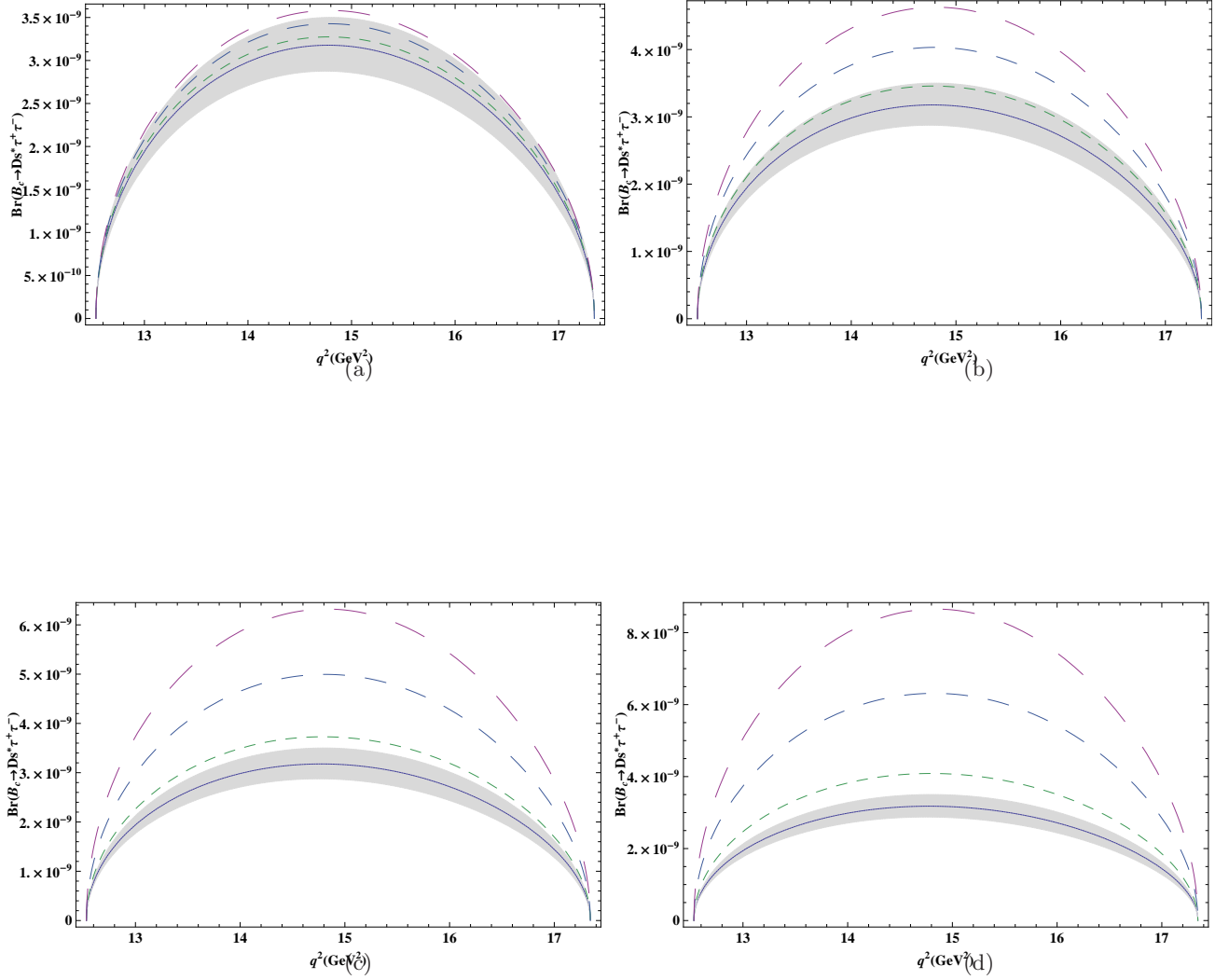


FIG. 2: The dependence of branching ratio of $B_c \rightarrow D_s^* \tau^+ \tau^-$ on q^2 for different values of $m_{t'}$ and $|V_{t'b}^* V_{t's}|$. The values of fourth generation parameters and the legends are same as in Fig.1

In Figs. 5 and 6 we have shown the dependence of longitudinal and transverse helicity fractions of D_s^* meson in $B_c \rightarrow D_s^* \mu^+ \mu^-$ decay on q^2 . We can establish from Fig. 5 (6) that the effects of the fourth generation on the longitudinal (transverse) helicity fractions of D_s^* are marked up in the $2.5 < q^2 \leq 13.5 \text{ GeV}^2$ region. Here one can see that the shift in the value of longitudinal and transverse helicity fractions from that of the SM value is small for the lower values of the fourth generation parameters ($m_{t'}$, $|V_{t'b}^* V_{t's}|$) which however becomes prominent when the values of these input parameters becomes maximum. The value of the longitudinal (transverse) helicity fraction of D_s^* meson decreases (increase) with the increment in the values of fourth generation input parameters $m_{t'}$ and $|V_{t'b}^* V_{t's}|$. *In contrast to analysis of same observable made in the UED model [41] where UED effects were mitigated by the WA contributions, we can see that in the present scenario the NP effects are very promising. Therefore, we expect that the experimental study of helicity fraction of D_s^* meson will serve as a handy tool to distinguish different NP scenarios.*

Performing the similar study when tau's are the final state leptons we have shown the dependence of longitudinal and transverse helicity fractions of the D_s^* meson in Figs. 7 and 8 against the square of the momentum transfer (q^2). Compared to the muon in the final state, the SM4 effects in this case are somewhat dim but still visible. Therefore,

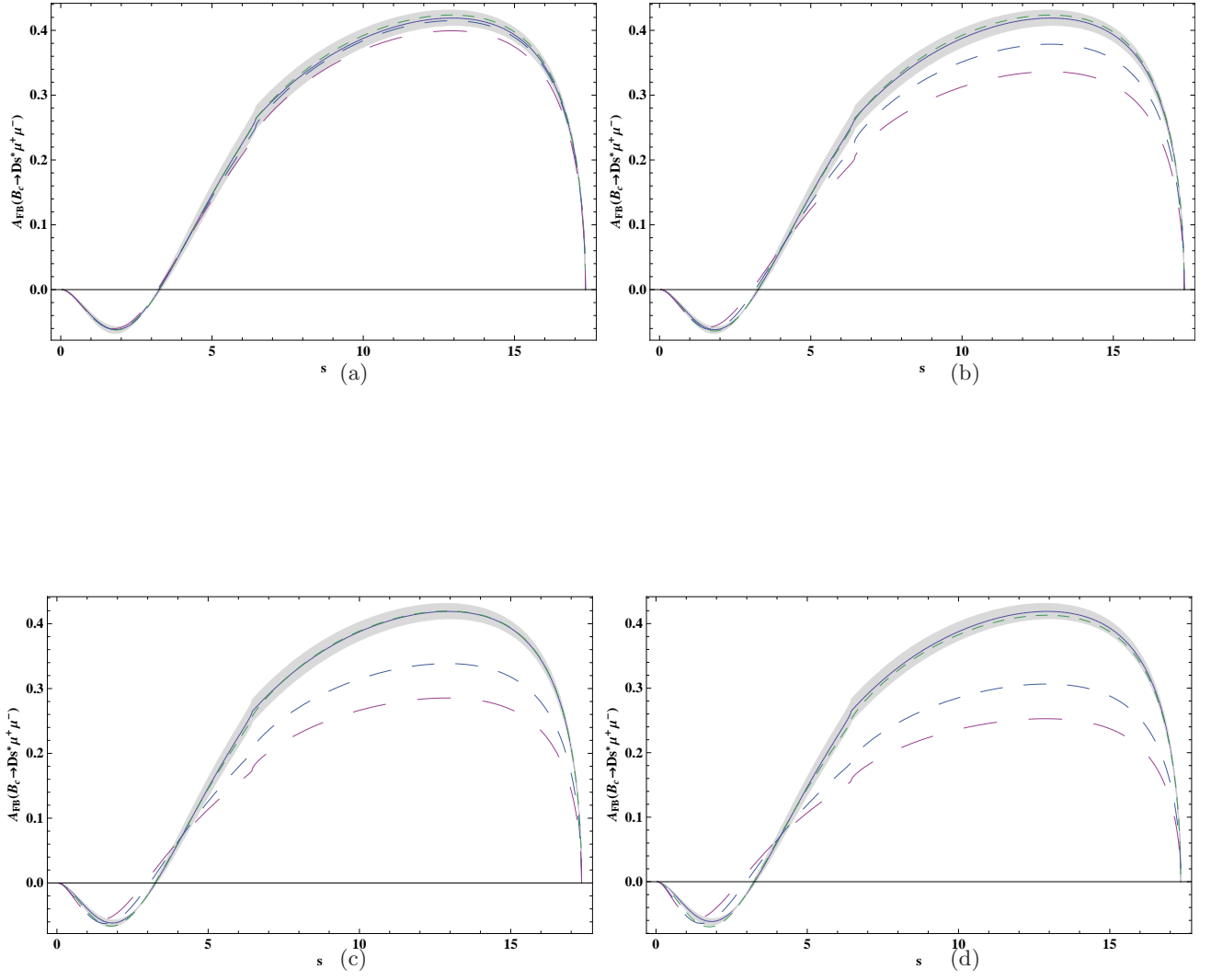


FIG. 3: The dependence of forward-backward asymmetry of $B_c \rightarrow D_s^* \mu^+ \mu^-$ on q^2 for different values of $m_{t'}$ and $|V_{t'b}^* V_{t's}|$. The values of fourth generation parameters and the legends are same as in Fig.1.

we expect that the experimental study of the helicity fraction will shed some light on the NP searches especially in $B_c \rightarrow D_s^* \mu^+ \mu^-$ decays.

Fig. 9 shows the dependence of longitudinal lepton polarization asymmetry for the $B_c \rightarrow D_s^* \mu^+ \mu^-$ decay on q^2 for different values of $m_{t'}$ and $|V_{t'b}^* V_{t's}|$. From Eq. (51) we can see that the WA contributions are canceled out and hence the effects of SM4 will be distinctively clear in longitudinal lepton polarization asymmetry which is depicted in Fig. 9. The value of longitudinal lepton polarization for muons is around -0.9 in the SM3 and we have significant deviation in this value in SM4. Just in the case of $m_{t'} = 600$ GeV and $|V_{t'b}^* V_{t's}| = 1.2 \times 10^{-2}$ the value of the longitudinal lepton polarization becomes -0.5 which will help us to see experimentally the SM4 effects in these decays. Similar effects can be seen for the final state tauons (c.f. Fig. 10). In this case the shift from the SM value is small compared to the muons in the final state because of the factor $\left(1 - \frac{4m_\tau^2}{q^2}\right)$ appear in the calculation of the longitudinal lepton polarization asymmetry.

The dependence of normal lepton polarization asymmetries for $B_c \rightarrow D_s^* \ell^+ \ell^-$ on the momentum transfer square are presented in Figs. 11 and 12. In terms of Eq. (52), it can be seen that it is proportional to the mass of the final

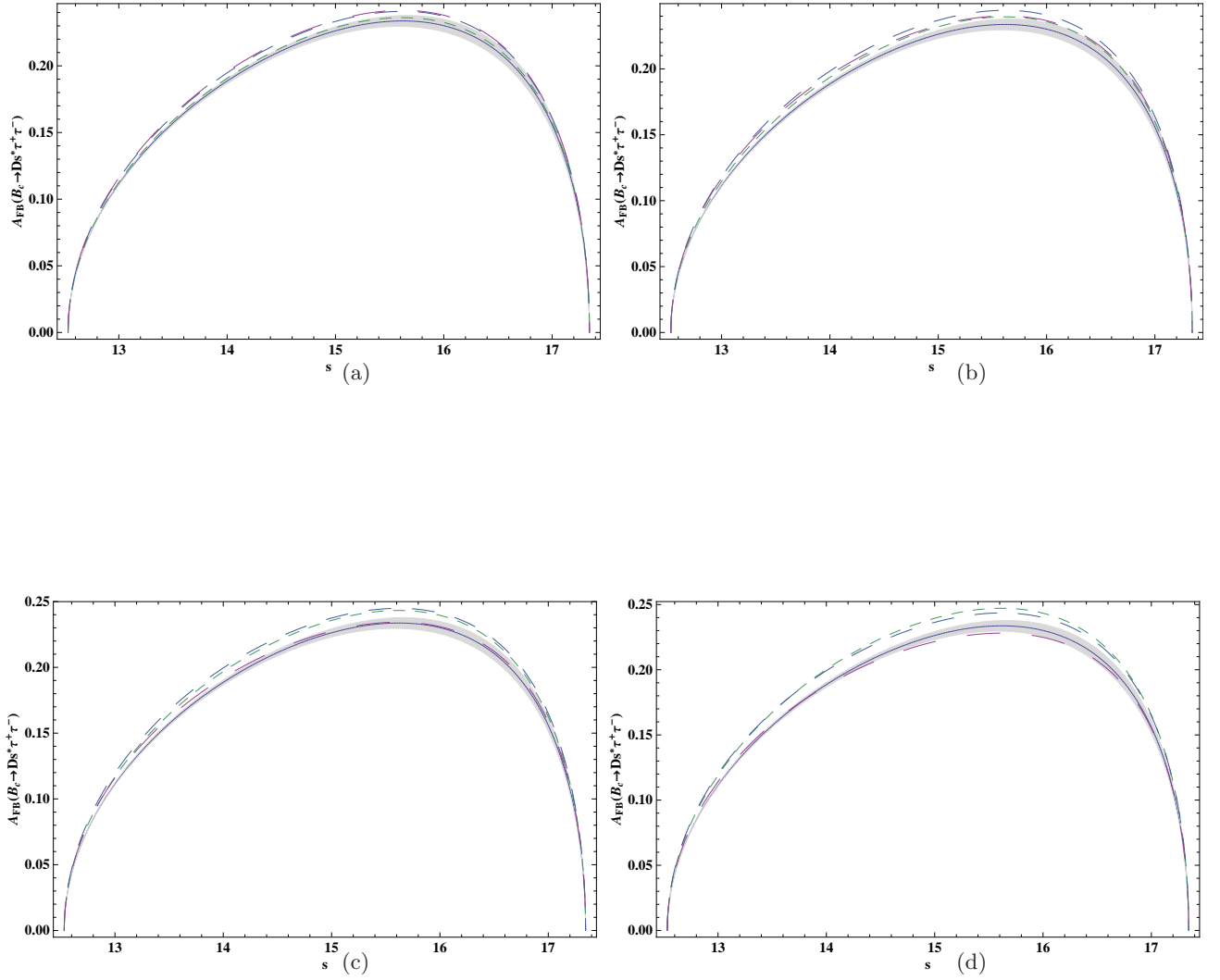


FIG. 4: The dependence of forward-backward asymmetry of $B_c \rightarrow D_s^* \tau^+ \tau^-$ on q^2 for different values of $m_{t'}$ and $|V_{t'b}^* V_{t's}|$. The values of fourth generation parameters and the legends are same as in Fig.1.

state lepton and for μ its values are expected to be small and Fig. 11 displays it in the SM as well as in the SM4 for the different values of fourth generation parameters. In SM4, there is a slight shift from the SM value which, however, due to its small value it is hard to measure experimentally. Now, for the $\tau^+ \tau^-$ channel, Eq. (52) we will have a large value of normal lepton polarization compared to the $\mu^+ \mu^-$ case in the SM. Fig. 12 shows that there is a significant decrease in the value of P_N in SM4 compared to the SM and its experimental measurement will give us some clear hints of the fourth generation of quarks.

Figs. 13 and 14 show the value of transverse lepton polarization both in the SM as well as in SM4 for $B_c \rightarrow D_s^* \ell^+ \ell^-$ decay. One can see that it is zero in the SM but non zero in sequential fourth generation SM (SM4). This non zero value comes from the interference of the Wilson coefficient for SM4 which are complex in SM4, see Eqs. (7, 53). The transverse lepton polarization is proportional to the lepton mass and as the imaginary part of the Wilson coefficients therefore, it is expected to be small which can also be seen from Figs. 13 and 14.

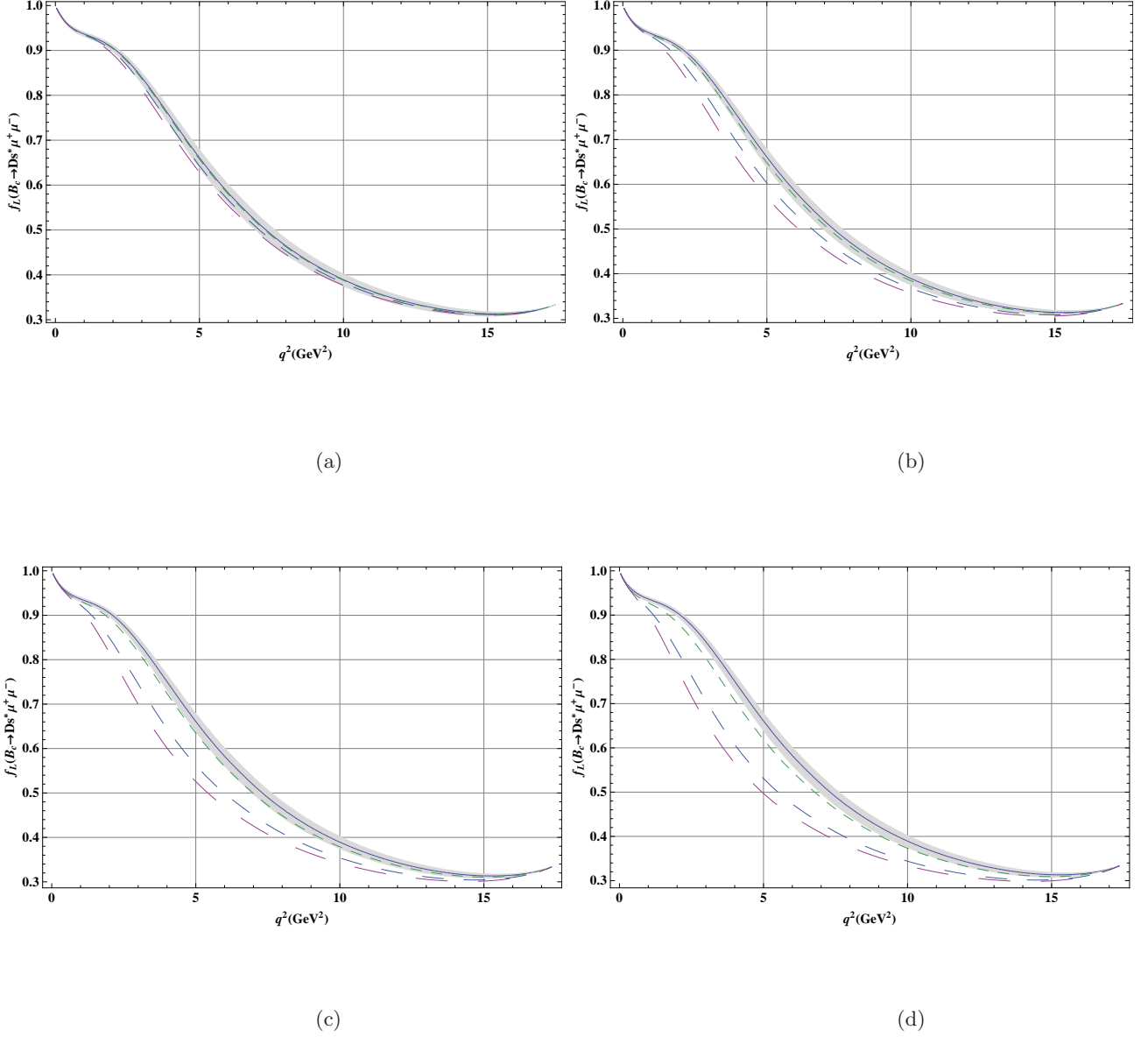


FIG. 5: The dependence of longitudinal helicity fractions of $B_c \rightarrow D_s^* \mu^+ \mu^-$ on q^2 for different values of $m_{t'}$ and $|V_{t'b}^* V_{t's}|$. The values of fourth generation parameters and the legends are same as in Fig.1.

V. SUMMARY AND CONCLUSION:

We have carried out the study of invariant mass spectrum, forward-backward asymmetries, polarization asymmetries of final state D_s^* meson and lepton in $B_c \rightarrow D_s^* \ell^+ \ell^-$ ($\ell = \mu, \tau$) decays in the Standard Model with extra sequential generation of quarks (SM4). Particularly, we analyze the effects of fourth generation up-type quark mass $m_{t'}$ and corresponding CKM matrix element $|V_{t'b}^* V_{t's}|$ to this process and our main outcomes can be summarized as follows:

- We have found that the branching ratios deviate sizably from that of the SM in almost all momentum transfer

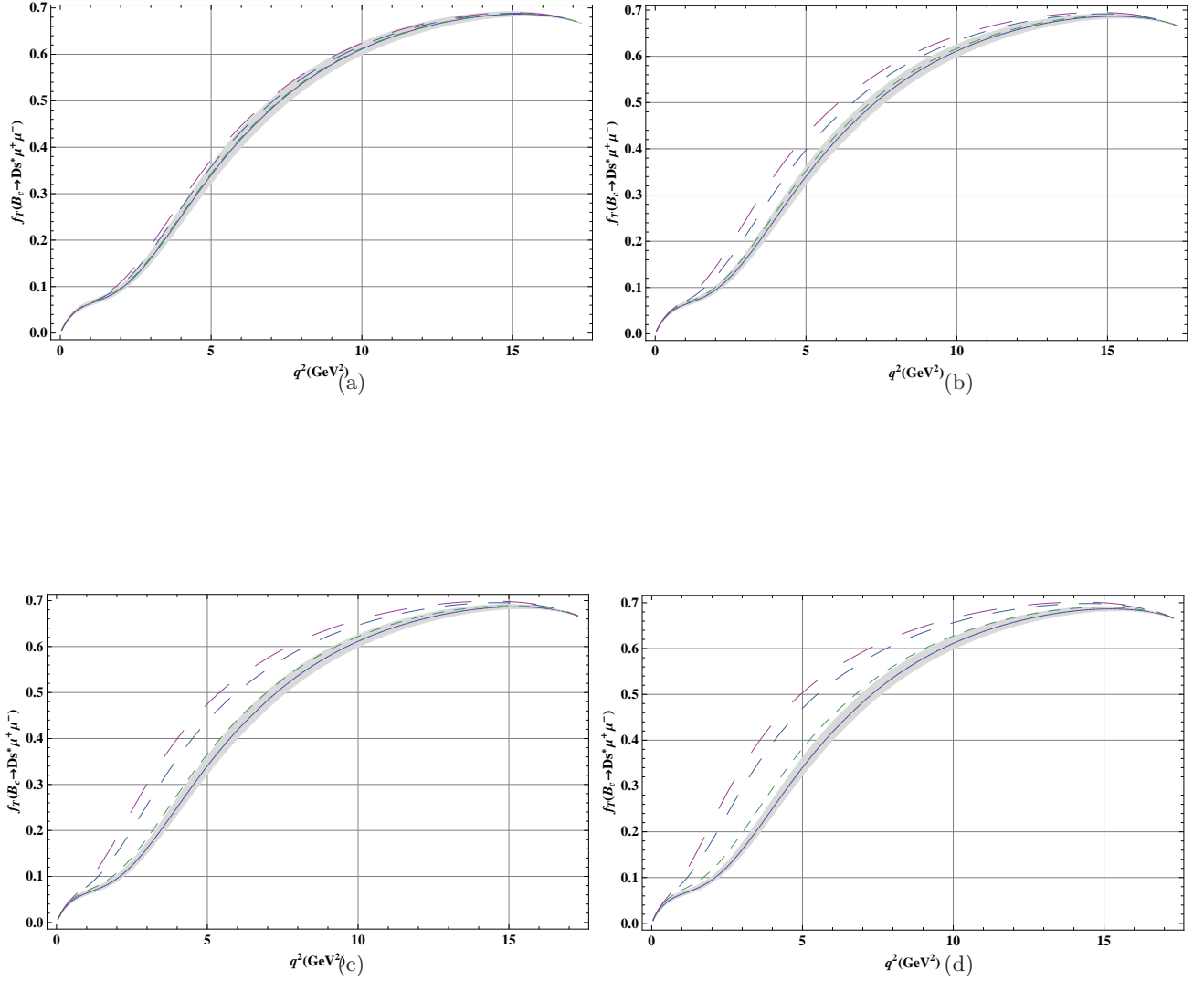


FIG. 6: The dependence of transverse helicity fractions of $B_c \rightarrow D_s^* \mu^+ \mu^-$ on q^2 for different values of $m_{t'}$ and $|V_{t'b}^* V_{t's}|$. The values of fourth generation parameters and the legends are same as in Fig.1.

region. The study shows that the \mathcal{BR} is an increasing function of the fourth generation parameters $m_{t'}$ and $V_{t'b} V_{t's}$. At maximum values of these parameters, i.e. $|V_{t'b}^* V_{t's}| = 0.012$ and $m_{t'} = 600$ GeV, the values of \mathcal{BR} increase approximately by 3 times of their SM values when the final leptons are muons or tauons. Hence the accurate measurement of the \mathcal{BR} for these decays is very important tool to explore physics beyond the SM.

- The value of the forward-backward asymmetry decreases significantly from that of the SM value in SM4 when the mass of the fourth generation quark varies from 300 GeV to 600 GeV. The value of the zero position of forward-backward asymmetry shifted towards the left for all values of $|V_{t'b}^* V_{t's}|$ in $B_c \rightarrow D_s^* \mu^+ \mu^-$ decay. This shifting is significant for large values of the fourth generation CKM matrix elements $|V_{t'b}^* V_{t's}|$ and fourth generation top quark mass $m_{t'}$. As it is almost free from the hadronic uncertainties therefore this shifting will help us to find clues of the SM4.
- The polarization effects of final state D_s^* meson and lepton are calculated in the sequential fourth generation SM4. It is found that the SM4 effects are very promising, which could be measured at present and future

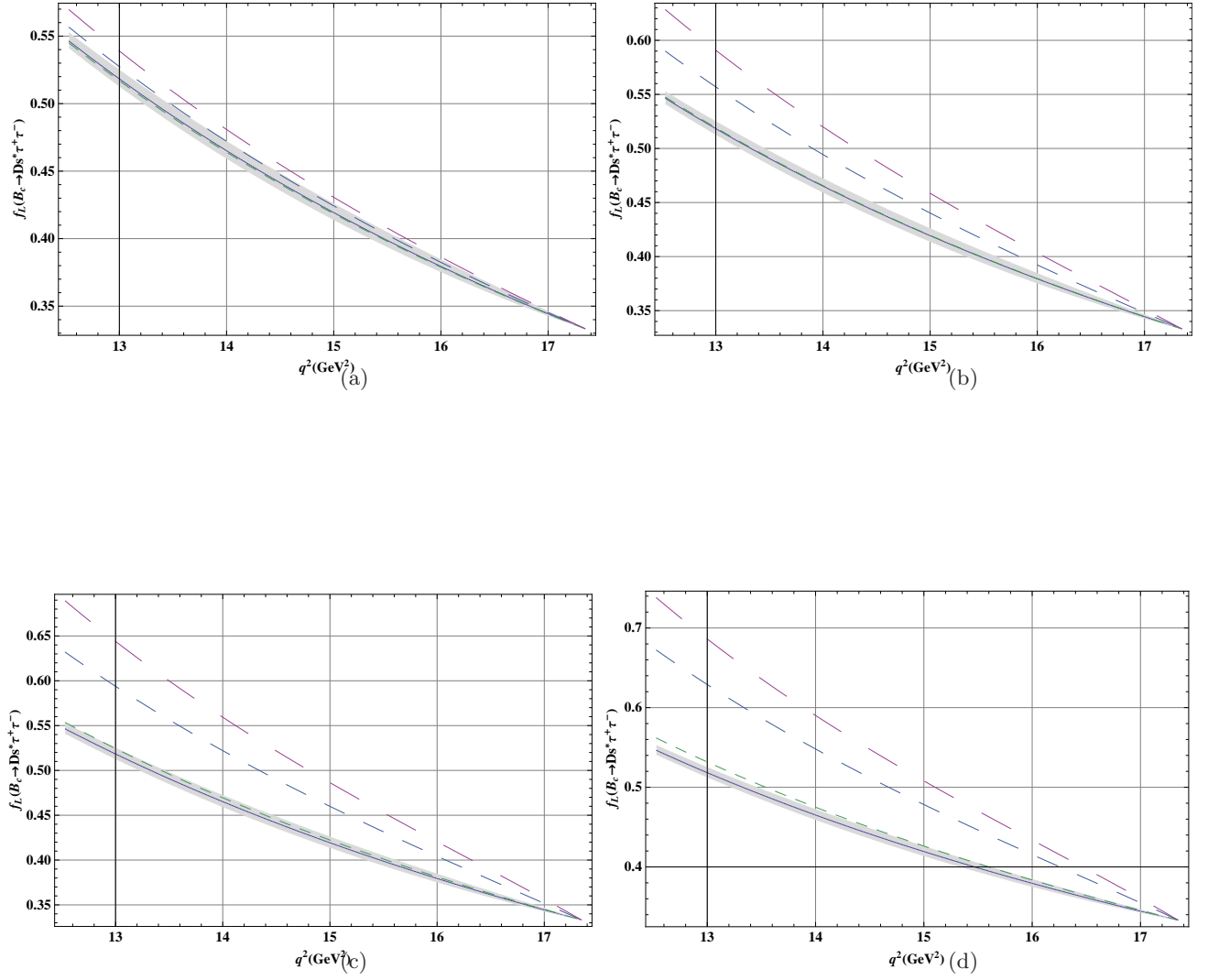


FIG. 7: The dependence of longitudinal helicity fractions of $B_c \rightarrow D_s^* \tau^+ \tau^-$ on q^2 for different values of $m_{t'}$ and $|V_{t'b}^* V_{t's}|$. The values of fourth generation parameters and the legends are same as in Fig.1.

experiments like LHCb where large numbers of $b\bar{b}$ pairs are expected to be produced.

In short, the precision measurements of these observables at Tevatron and LHC will help us to find the indications of new physics encoded in the fourth generation parameters such as $V_{t'b}^* V_{t's}$ and $m_{t'}$.

Acknowledgements

Helpful discussions with Prof. Riazuddin and Prof. Fayyazuddin are greatly acknowledged. M. J. A acknowledge the grant provided by Quaid-i-Azam University from University Research Funds.

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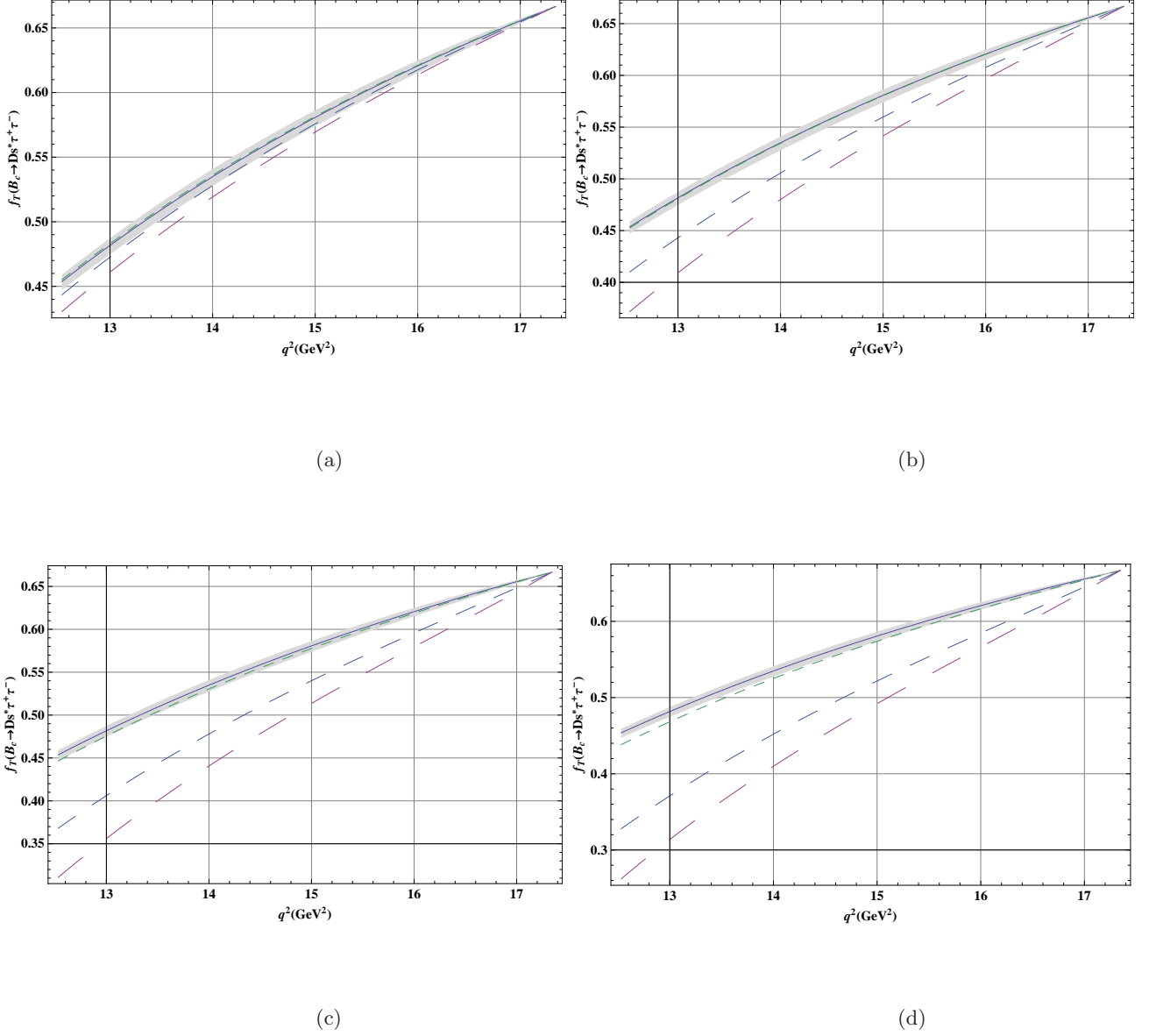


FIG. 8: The dependence of transverse helicity fractions of $B_c \rightarrow D_s^* \tau^+ \tau^-$ on q^2 for different values of $m_{t'}$ and $|V_{t'b}^* V_{t's}|$. The values of fourth generation parameters and the legends are same as in Fig.1.

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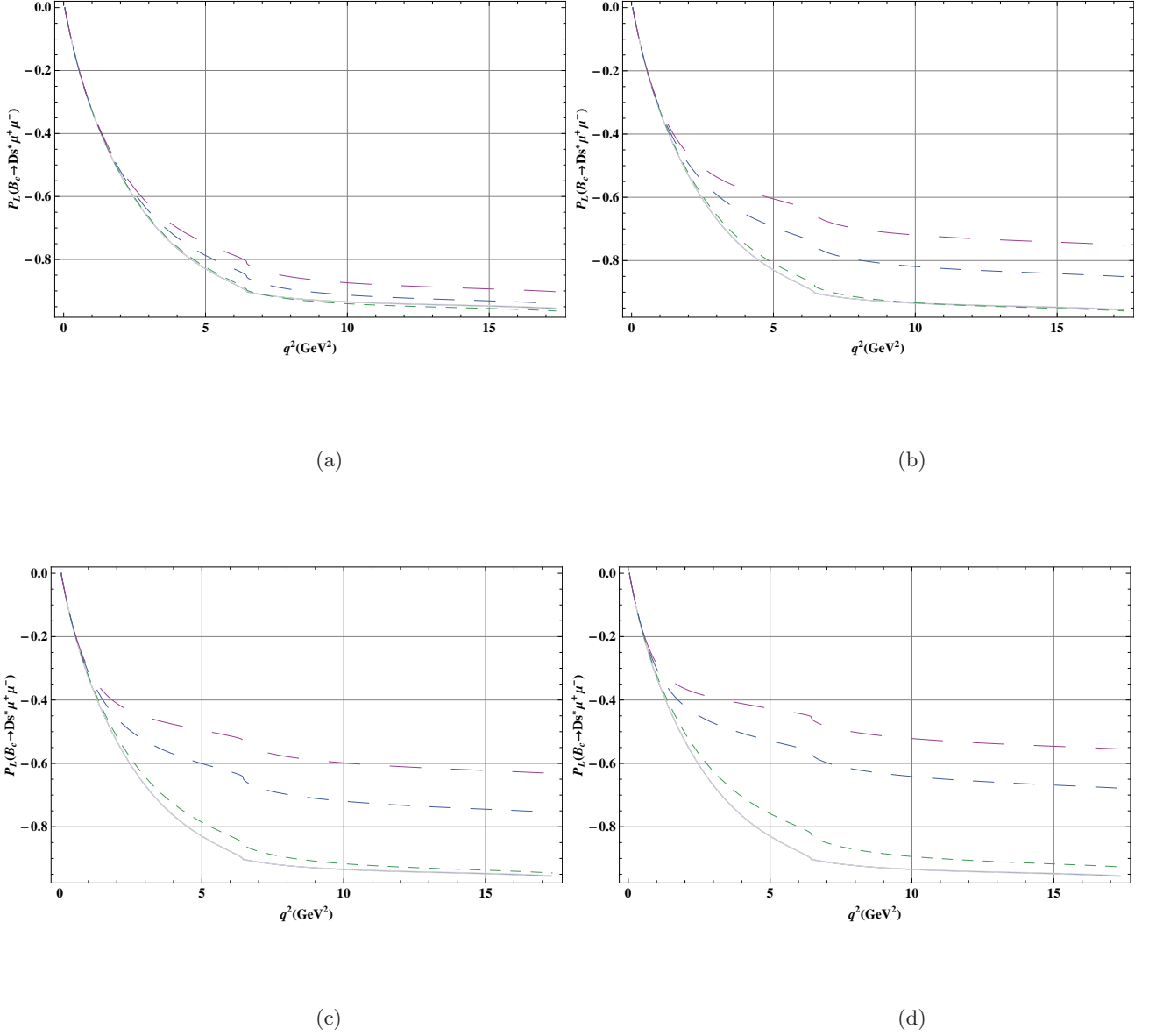


FIG. 9: The dependence of longitudinal lepton polarization of $B_c \rightarrow D_s^* \mu^+ \mu^-$ on q^2 for different values of $m_{\nu'}$ and $|V_{t'b}^* V_{t's}|$. The values of fourth generation parameters and the legends are same as in Fig.1.

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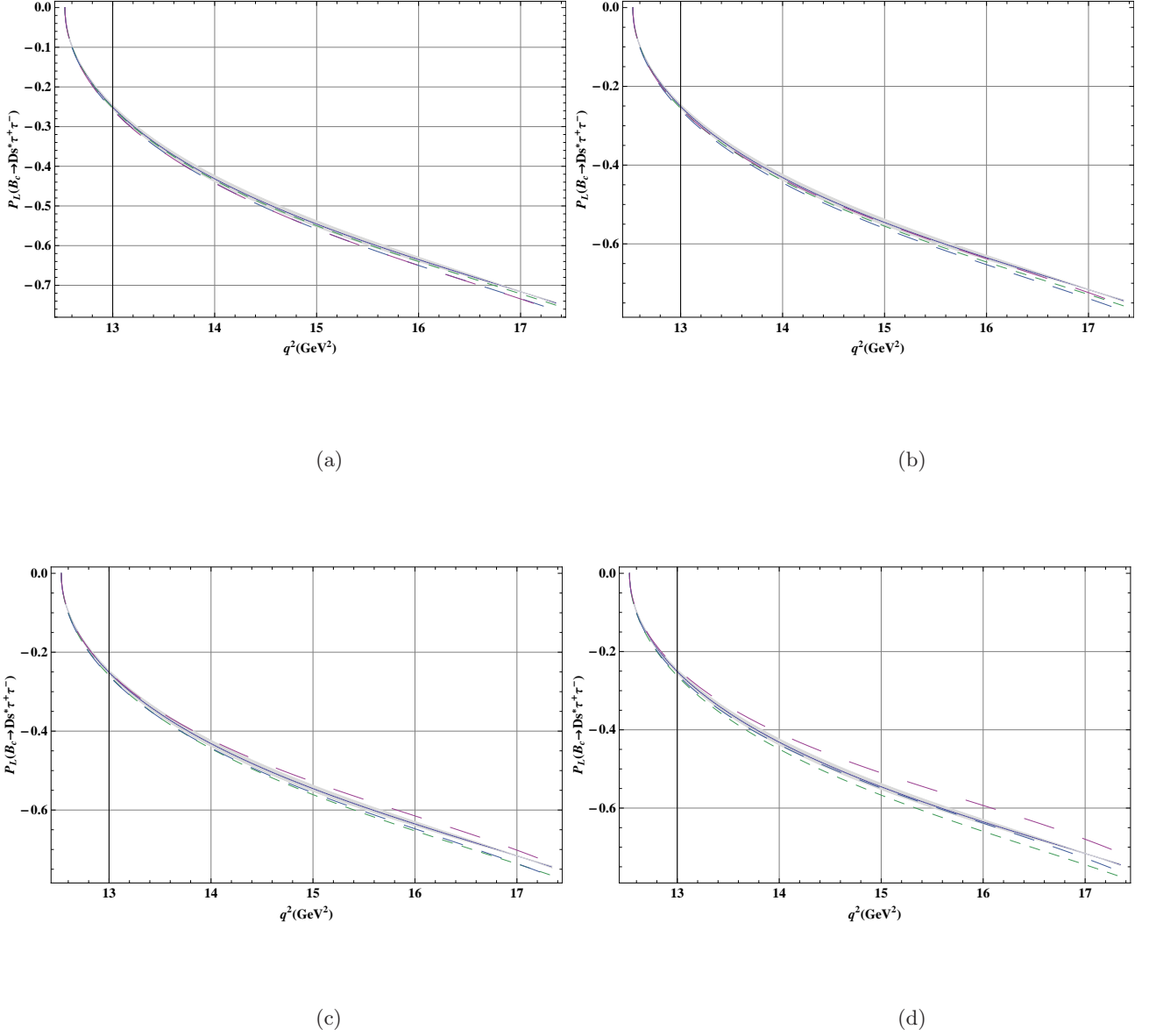


FIG. 10: The dependence of longitudinal lepton polarization of $B_c \rightarrow D_s^* \tau^+ \tau^-$ on q^2 for different values of $m_{l'}$ and $|V_{l'b}^* V_{l's}|$. The values of fourth generation parameters and the legends are same as in Fig.1.

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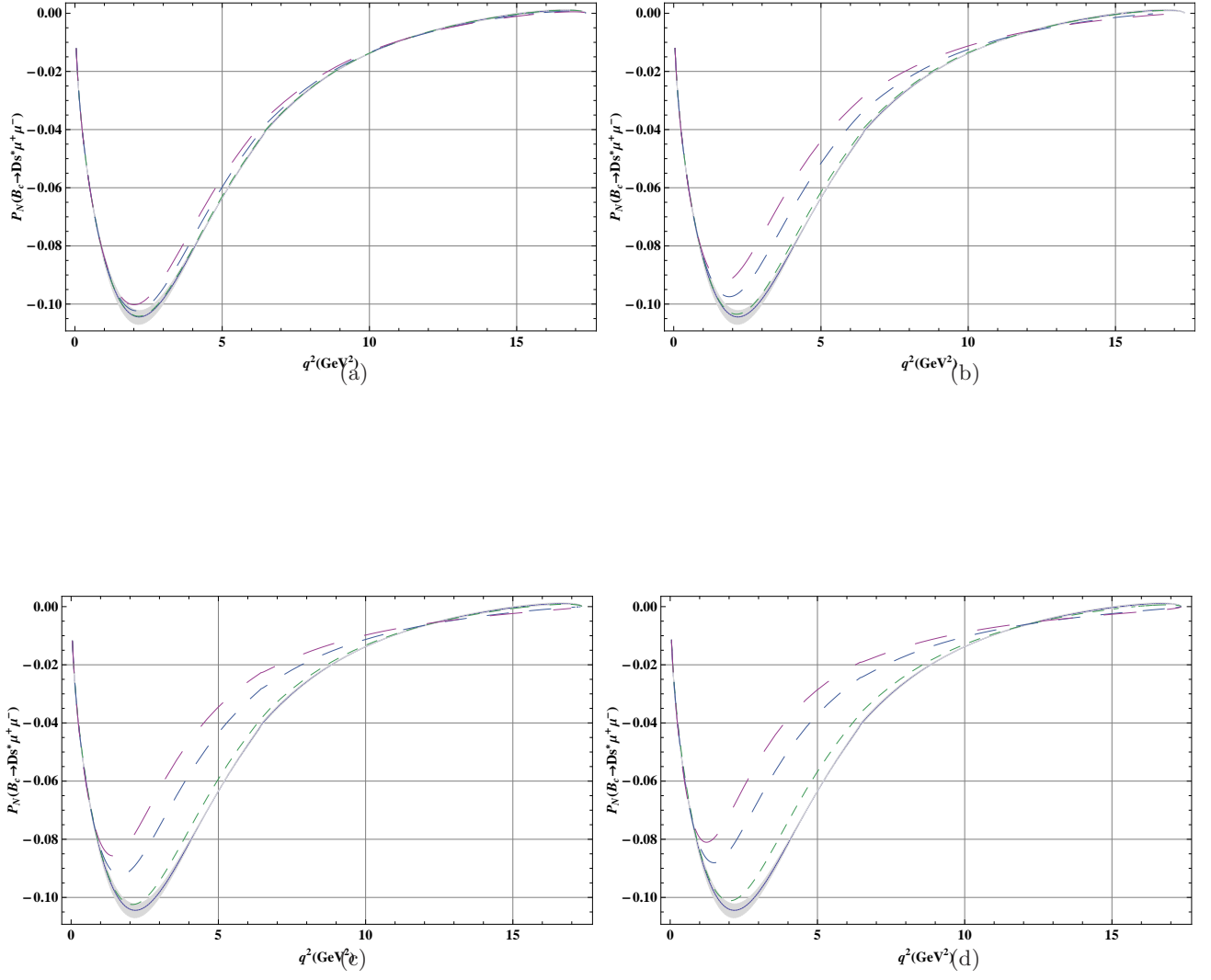


FIG. 11: The dependence of normal lepton polarization of $B_c \rightarrow D_s^* \mu^+ \mu^-$ on q^2 for different values of $m_{t'}$ and $|V_{t'b}^* V_{t's}|$. The values of fourth generation parameters and the legends are same as in Fig.1.

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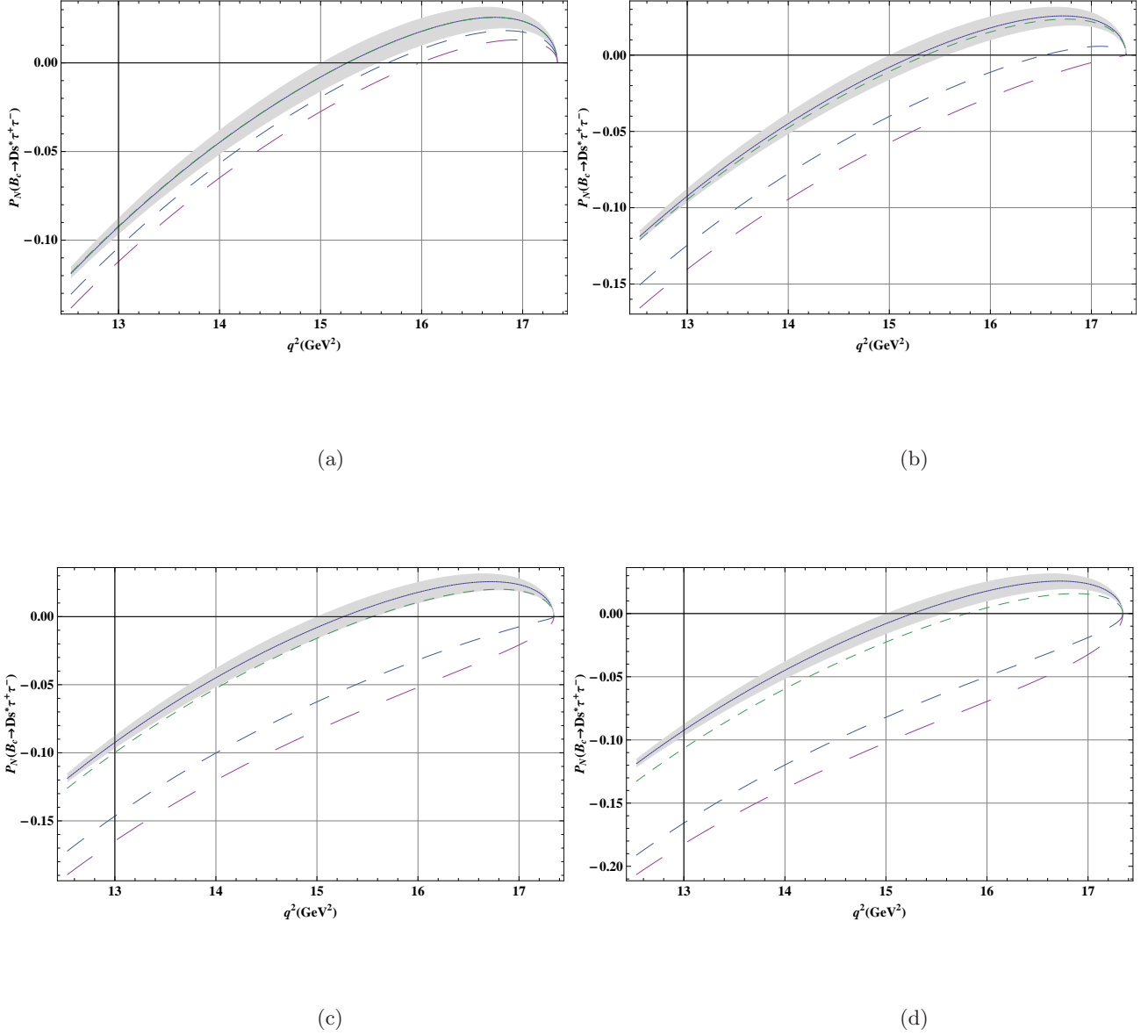


FIG. 12: The dependence of normal lepton polarization of $B_c \rightarrow D_s^* \tau^+ \tau^-$ on q^2 for different values of $m_{t'}$ and $|V_{t'b}^* V_{t's}|$. The values of fourth generation parameters and the legends are same as in Fig.1.

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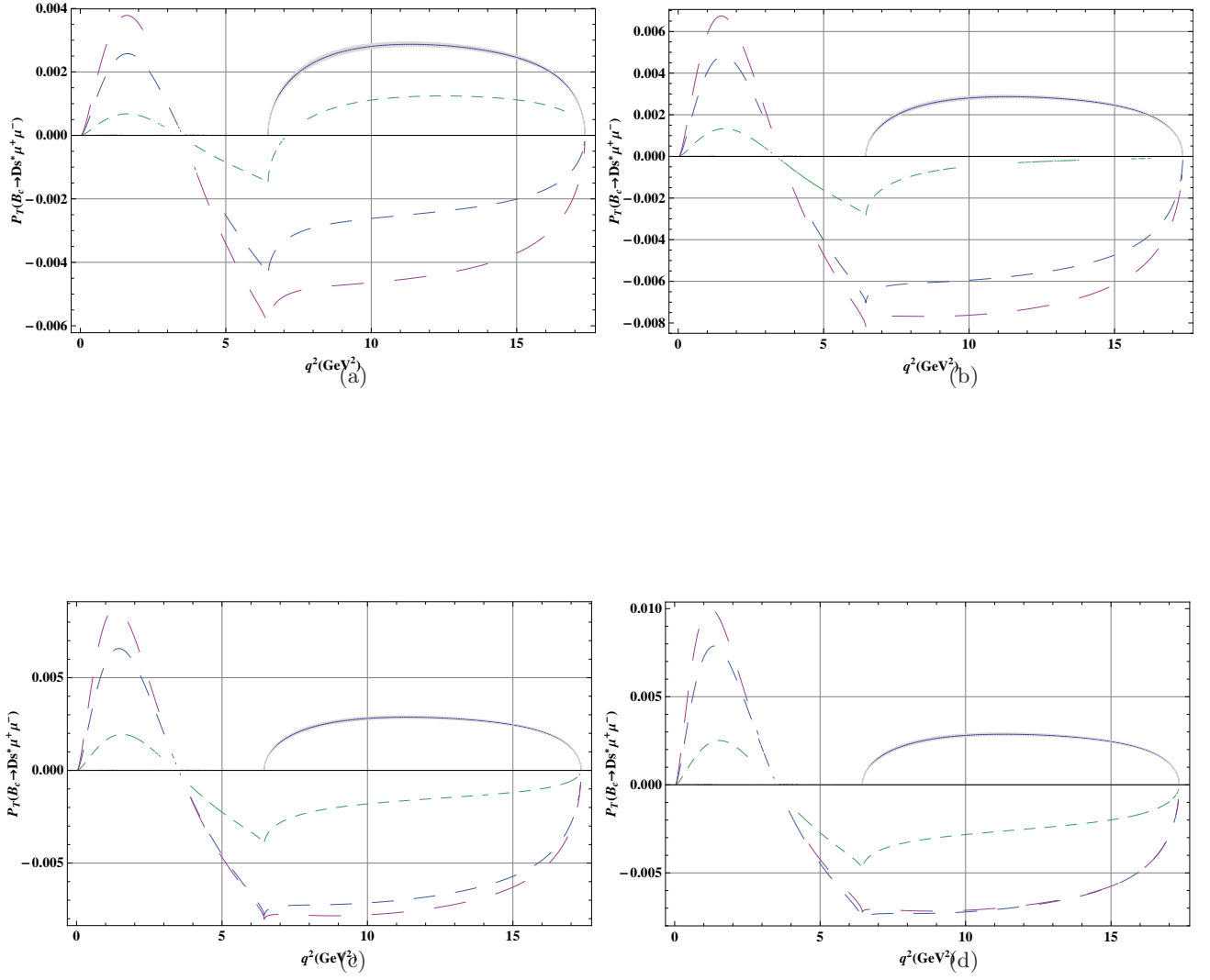


FIG. 13: The dependence of transverse lepton polarization of $B_c \rightarrow D_s^* \mu^+ \mu^-$ on q^2 for different values of $m_{t'}$ and $|V_{t'b}^* V_{t's}|$. The values of fourth generation parameters and the legends are same as in Fig.1.

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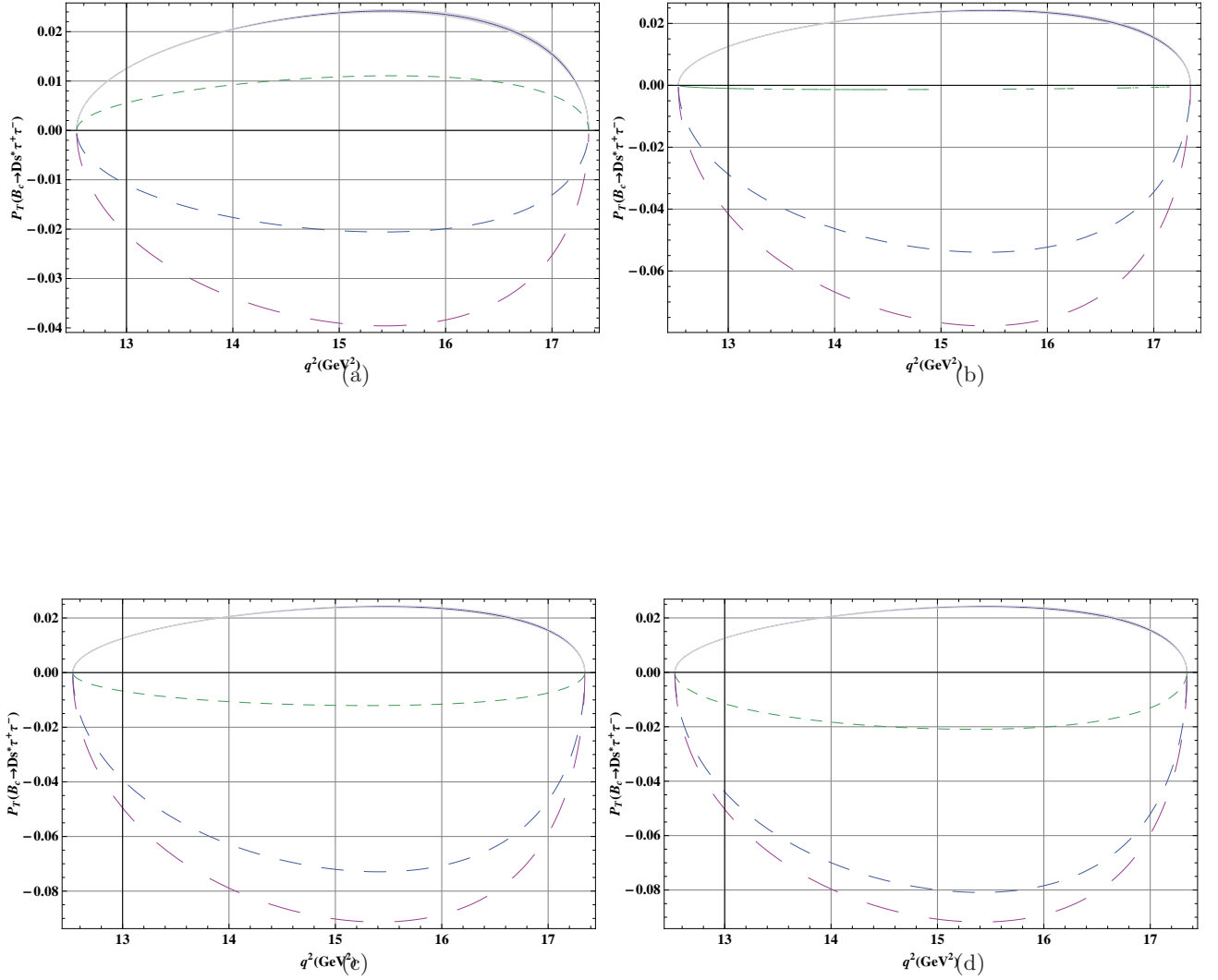


FIG. 14: The dependence of transverse lepton polarization of $B_c \rightarrow D_s^* \tau^+ \tau^-$ on q^2 for different values of $m_{t'}$ and $|V_{t'b}^* V_{t's}|$. The values of fourth generation parameters and the legends are same as in Fig.1.

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